Boundary-Layer Theory, Strong-Coupling Series, and Large-Order Behavior

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Abstract
The introduction of a lattice converts a singular boundary-layer problem in the continuum into a regular perturbation problem. However, the continuum limit of the discrete problem is extremely nontrivial and is not completely understood. This paper examines two singular boundary-layer problems taken from mathematical physics, the instanton problem and the Blasius equation, and in each case examines two strategies, Padé resummation and variational perturbation theory, to recover the solution to the continuum problem from the solution to the associated discrete problem. Both resummation procedures produce good and interesting results for the two cases, but the results still deviate from the exact solutions. To understand the discrepancy a comprehensive large-order behavior analysis of the strong-coupling lattice expansions for each of the two problems is done.

I Introduction

In this paper we report some major advances in understanding (albeit not a complete solution to) a difficult general class of problems in mathematical physics. We consider here the conversion of a continuum problem into a discrete problem by the insertion of a lattice spacing parameter $a$, the solution of the continuum problem on the lattice, and the subsequent extremely subtle continuum limit $a \to 0$.

Almost every continuum physics problem is singular as a function of the parameters in the problem. As a result, only rarely does the perturbation series take the form of a Taylor series having a nonzero radius of convergence. As an elementary example, consider the algebraic polynomial equation

$$\epsilon x^3 + x - 1 = 0.$$  (1)
This problem is singular in the limit $\epsilon \to 0$. In this limit, the degree of the polynomial changes from three to one and thus two of the roots abruptly disappear. As a consequence, a perturbative solution to this problem [expressing the roots $x(\epsilon)$ as series in powers of $\epsilon$] yields expressions that are more complicated than Taylor series.

A more elaborate example of a singular problem is the time-independent Schrödinger equation

$$-\frac{\hbar^2}{2M} \nabla^2 \Psi(x) + [V(x) - E] \Psi(x) = 0. \quad (2)$$

In the classical limit $\hbar \to 0$ this differential equation abruptly becomes an algebraic equation, and thus the general solution no longer contains any arbitrary constants or functions and, as a result, it can no longer satisfy the initial conditions. We know that for small $\hbar$ the solution is not Taylor-like but rather is a singular exponential in WKB form:

$$\Psi(x) \sim e^{S(x)/\hbar} \quad (\hbar \to 0). \quad (3)$$

In the study of quantum field theory, it is well known that infinities appear in the perturbative expansion in powers of the coupling constant. There are two kinds of infinities. The first kind, which is due to the point-like nature of the interaction, requires the use of renormalization. The second kind, which is due to singularities in the complex-coupling-constant plane, forces the perturbation series to have a zero radius of convergence.

A quantum field theory can be regulated by introducing a lattice spacing. The resulting discrete theory is completely finite and can be studied numerically by using various kinds of numerical methods such as Monte Carlo integration. However, the underlying singular nature of the continuum quantum field theory resurfaces in the continuum limit $a \to 0$. The introduction of a lattice spacing and the singular nature of the continuum limit was investigated in a series of papers by Bender et al. [1–9].

A quantum field theory is just one instance in which discretization regulates and eliminates the singular nature of the problem. It is also known that introducing a lattice spacing converts a boundary-layer problem, which is a singular perturbation problem, into a regular perturbation problem [10–12]. A boundary-layer problem is a differential-equation-boundary-value problem in which the highest derivative of the differential equation is multiplied by a small parameter $\epsilon$. Consider as an example

$$\epsilon y''(x) + a(x)y'(x) + b(x)y(x) = c(x), \quad (4)$$

where the boundary conditions on the function $y(x)$ typically have a form such as

$$y(0) = A, \quad y(1) = B. \quad (5)$$

This boundary-value problem is singular because in the limit $\epsilon \to 0$ one of the solutions abruptly disappears and the limiting solution is not able to satisfy the two boundary conditions in (5). The usual way to solve the boundary-value problem (3) – (5) is to decompose the interval $0 \leq x \leq 1$ into two regions, an outer region, in which the solution varies slowly as a function of $x$, and an inner region or boundary-layer region, in which the solution varies rapidly as a function of $x$. The boundary-layer region is a narrow region whose thickness is typically of order $\epsilon$ or some power of $\epsilon$ [13].

An important example of a boundary-layer problem is the instanton equation

$$\epsilon^2 f''(x) + f(x) - f^3(x) = 0, \quad (6)$$

with the associated boundary conditions

$$f(0) = 0, \quad f(\infty) = 1. \quad (7)$$
The exact solution to this instanton problem is
\[ f(x) = \tanh \frac{x}{\epsilon \sqrt{2}}. \tag{8} \]

Note that the solution \( f(x) \) varies rapidly at the origin \( x = 0 \) over a region of thickness \( \epsilon \); this is the boundary-layer region. The solution varies slowly (it is approximately 1) outside of this region. The outer region consists of those \( x \) not near the origin.

A novel way to solve the instanton problem is to discretize it by introducing a lattice. On the lattice, the differential equation becomes a difference equation that can easily be solved perturbatively. In the continuum limit, as the lattice spacing vanishes, we then obtain a strong-coupling expansion that must be evaluated by means of a Padé or a variational perturbation theory method. To illustrate the approach our objective will be to calculate the slope of the instanton at \( x = 0 \), which from (8) has the value
\[ f'(0) = \frac{1}{\epsilon \sqrt{2}}. \tag{9} \]

We introduce a lattice with lattice spacing \( a \) so that the real axis is discretized in steps of width \( a \). The spatial coordinate reads \( x_n = na \), where the function \( f(x) \) assumes the value \( f_n = f(x_n) \). On the lattice the second spatial derivative in (6) becomes
\[ f''(x) \to \frac{f_{n+1} - 2f_n + f_{n-1}}{a^2}. \tag{10} \]

Thus, from the instanton equation (6) we obtain the difference equation
\[ \frac{\epsilon^2}{a^2} (f_{n+1} - 2f_n + f_{n-1}) + f_n - f_n^3 = 0, \tag{11} \]
where the boundary values follow from (7):
\[ f_0 = 0, \quad f_\infty = 1. \tag{12} \]

The natural expansion parameter now is \( \epsilon^2/a^2 \), to which we assign the name \( \delta \):
\[ \delta \equiv \frac{\epsilon^2}{a^2}. \tag{13} \]

The singular perturbation problem in the continuum [whose solution \( f(x) \) in (8) does not possess a Taylor expansion in powers of \( \epsilon \), has become a regular perturbation problem. That is, we can now expand the solution \( f_n \) to the difference equation (11) as a Taylor series in powers of \( \delta \):
\[ f_n = a_{n,0} + a_{n,1}\delta + a_{n,2}\delta^2 + \ldots. \tag{14} \]

We impose the boundary values (12) by requiring that
\[ a_{0,0} \equiv 0 \quad \text{and} \quad a_{n,0} \equiv 1 \quad (n \geq 1). \tag{15} \]

Inserting the ansatz (14) into the difference equation (11), we get the recursion relation (10)
\[ a_{n,j} = \frac{1}{2}a_{n+1,j-1} + a_{n,j-1} + \frac{1}{2}a_{n-1,j-1} - \sum_{k=1}^{j-1} a_{n,k}a_{n,j-k} - \frac{1}{2} \sum_{k=1}^{j-1} \sum_{l=1}^{j-k} a_{n,k}a_{n,l}a_{n,j-k-l}. \tag{16} \]
For the first derivative at the origin $x = 0$ this leads to the series

$$f'(0) = \lim_{a \to 0} \frac{f_1 - f_0}{a} = \lim_{a \to 0} \frac{f_1}{a} = \lim_{a \to 0} \frac{1}{a} \sum_{j=0}^{\infty} a_{1,j} \delta^j = \lim_{a \to 0} \frac{1}{a} \left( 1 - \frac{\delta}{2} + \frac{\delta^2}{8} + \frac{11\delta^4}{128} + \ldots \right).$$

(17)

We have calculated the coefficients $a_{n,j}$ with the help of Maple V R7 up to order $j = 200$. The first 20 numbers are given in Table I. A complete list of these coefficients can be found on the webpage of the author FW [14]. Note that the expansion parameter $\delta$ in (17) is not small but rather tends to infinity in the limit as the lattice spacing $a$ approaches zero. Using the parameter $\delta$ defined in (13) we rewrite the series (17) as

$$f'(0) = \frac{1}{\epsilon} \lim_{\delta \to \infty} \sqrt{\delta} \left( 1 - \frac{\delta}{2} + \frac{\delta^2}{8} + \frac{11\delta^4}{128} + \ldots \right).$$

(18)

Taking into account the exact result (9), we obtain the identity

$$\frac{1}{\sqrt{2}} = \lim_{\delta \to \infty} \sqrt{\delta} \left( 1 - \frac{\delta}{2} + \frac{\delta^2}{8} + \frac{11\delta^4}{128} + \ldots \right).$$

(19)

The purpose of this paper is to examine equations like (19). This equation shows that the singular nature of the instanton problem has resurfaced in the continuum limit $\delta \to \infty$ of the lattice expansion. The expression on the right side of (19) should have the value $1/\sqrt{2} = 0.7071067812\ldots$, but it is not at all obvious why this is so, and the objective of this paper is to analyze this difficult and subtle limit.

This paper is organized as follows. In Sec. II we use Padé techniques to perform the limit in (19). We will see that while the results are not bad (the accuracy is about 1%), better methods are needed. We perform the Padé analysis to much higher order than has ever been done before and we discover a new qualitative behavior that has not yet been observed. In Sec. III we try the use of the variational perturbation theory techniques introduced by Kleinert to perform the sum in (19). These techniques increase the accuracy by a factor of about 10, but they still do not give the exact result. While variational perturbation theory works very well in summing the strong-coupling series for the ground-state energy of the anharmonic oscillator [15], and for the critical exponents of second-order phase transitions [16], we show that the series in (19) is at the very edge of validity for Kleinert’s methods. We then examine the large-order behavior of the terms of the sum in (19) in Sec. IV. We show definitively that the Taylor expansion has a nonzero radius of convergence and thus, on the lattice, the instanton problem is a regular perturbation problem. In Sec. V we turn to a more difficult singular perturbation problem; namely, the Blasius equation of fluid dynamics. We use the same approach as for the instanton equation. In Secs. VI, VII, and VIII we study the summation of the lattice perturbation expansion using Padé and variational methods and we examine the large-order behavior of the lattice perturbation series. We find that Padé methods give good but not excellent results and that variational perturbation theory is better than Padé. Again, the series we need to evaluate in the continuum limit lies at the very edge of validity for Kleinert’s methods. We also find that, unlike the lattice perturbation expansion coefficients for the instanton problem, the sign pattern of the Blasius weak-coupling series does not alternate. Rather, it is governed by a cosine function with a frequency different from $\pi$. 

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II Padé Resummation for the Instanton Equation

In this section we examine what happens if we attempt to evaluate the right side of (19) by using Padé techniques. Padé resummation has already been applied to the instanton problem up to 50th order [10]. However, we have been able to perform the procedures to much higher orders. We have discovered that remarkable and unsuspected new phenomena occur just a few orders beyond what has been computed.

The procedure is as follows. Consider the formal Frobenius series

\[ S(\delta) = \delta^M \sum_{n=0}^{\infty} a_n \delta^n, \tag{20} \]

where \( M \) is a non-negative number. Raising this series to the power \( 1/M \), inverting the right hand side and re-expanding, we obtain

\[ S^{1/M}(\delta) = \frac{\delta}{\sum_{n=0}^{\infty} b_n \delta^n}, \tag{21} \]

with new expansion coefficients \( b_n \). Assuming we know the first \( N + 1 \) terms of the original power series in (20), we raise equation (21) to the power \( N \). We then truncate the summation at \( n = N \), finally getting

\[ S^{N/M}(\delta) = \frac{\delta^N}{\sum_{n=0}^{N} c_n^{(N)} \delta^n}, \tag{22} \]

where we have re-expanded and obtained new expansion coefficients \( c_n^{(N)} \). In the limit \( \delta \to \infty \), only the \( N \)th term in the denominator survives and we obtain the approximant

\[ (S_N)^{N/M} \equiv \lim_{\delta \to \infty} S^{N/M}(\delta) = \lim_{\delta \to \infty} \frac{\delta^N}{\sum_{n=0}^{N} c_n^{(N)} \delta^n} = \frac{1}{c_n^{(N)}}, \tag{23} \]

The approximant \( S_N = \left( c_n^{(N)} \right)^{-M/N} \) is the zeroth-order survivor of the limiting process. Also, taking into account the first-order correction we observe that, as in the case of variational perturbation theory (see Sec. III), there is an approach to scaling. In the limit \( \delta \to \infty \) the Frobenius series \( S(\delta) \) in Eq. (20) converges to a constant \( C \). Additionally, the approach to scaling, following from the Padé resummation (23), reveals how fast it converges:

\[ S(\delta) \sim C + C' \delta^{-1} \quad (\delta \to \infty). \tag{24} \]

We now apply this procedure to the boundary-layer problem [11]. Recall that the weak-coupling coefficients for the first 20 coefficients \( a_{1,j} \) obtained from (16) are shown in Table 1 and that more can be found in [14]. Resumming the series (14) for \( n = 1 \),

\[ f_1 = \sum_{j=0}^{N} a_{1,j} \delta^j, \tag{25} \]

according to the Padé procedure (23) with \( M = 1/2 \) as follows from (19) and evaluating the approximants \( S_N = \left( c_n^{(N)} \right)^{-M/N} \), we get the numbers listed in Table 2.
Compared with the numerical solution $1/\sqrt{2} \approx 0.7171067812$, this strong-coupling expansion seems to converge quite well. However, when we go to higher orders, we find that the numbers drop below the exact solution and assume a minimum at $N = 24$, where the approximant has the value $S_{24} \approx 0.70198319$. The approximants then rise again, cross the exact solution at $N = 41$ and become complex at $N = 52$. The appearance of complex numbers is a consequence of taking the $N$th root in equation (23) when the coefficients $c_N^{(N)}$ become negative. This phenomenon has not been observed before in the course of using this Padé procedure. The imaginary part then becomes smaller and smaller as $N$ rises. Abruptly, at $N = 68$, the approximants become real again. As one can see from the spikes in Fig. 1 this pattern is repeated for higher $N$. Note that the figure only shows the real part of the Padé approximant $S_N$.

Apparently, the sequence of approximants $S_N$ does not converge. The singular nature of the instanton equation has the effect of making the Padé approximants behave like the partial sums of a divergent (asymptotic) series; at first the partial sums appear to converge to a limit, and then they veer off. In the case of the Padé’s shown in Fig. 1 the approximants approach to within 1% of the correct limit before veering off. It appears that another more powerful resummation technique is needed to treat the expression in (19). In the next section we apply a technique due to Kleinert.

### III Variational Perturbation Theory for the Instanton Equation

Kleinert has developed a technique in the context of the ground-state energy of the anharmonic oscillator [15] and of critical exponents of second-order phase transitions [16] for summing divergent perturbation series. This technique, known as Kleinert’s square-root trick, is described below.

Consider a weak-coupling series

$$f_N(\delta) = \sum_{n=0}^{N} f_n \delta^n,$$

which is truncated at order $N$. Rewrite this weak-coupling expansion by introducing an auxiliary scaling parameter $\kappa$ [15, 16]:

$$f_N(\delta) = \kappa^p \sum_{n=0}^{N} f_n \left( \frac{\delta}{\kappa^q} \right)^n \bigg|_{\kappa=1},$$

which is set to $\kappa = 1$ later. The square-root trick now reads

$$\kappa \to \sqrt{K^2 + \kappa^2 - K^2} = K \sqrt{1 + \delta r},$$

where $K$ is a “dummy” scaling parameter and

$$r = \frac{1}{\delta} \left( \frac{\kappa^2}{K^2} - 1 \right).$$

In the case of the anharmonic oscillator, $K$ is the frequency $\Omega$ of a trial harmonic oscillator [15].

Substituting (28) into the truncated weak-coupling series (27), we obtain

$$f_N(\delta, K) = \sum_{n=0}^{N} f_n K^{p-nq} (1 + \delta r)^{(p-nq)/2} \delta^n.$$
The factor \((1 + \gamma r)^\alpha\) with \(\alpha \equiv (p - nq)/2\) can be expanded by means of generalized binomials according to
\[
(1 + \gamma r)^\alpha = \sum_{k=0}^{N-n} \binom{\alpha}{k} (\gamma r)^k \delta^n = \sum_{k=0}^{N-n} \binom{\alpha}{k} \left( \frac{1}{K^2} - 1 \right)^k \delta^n, 
\] (31)
where we have used (28) and finally have set \(\kappa \equiv 1\). The binomial is defined as
\[
\binom{\alpha}{k} = \frac{\Gamma(\alpha + 1)}{\Gamma(k + 1)\Gamma(\alpha + k + 1)}. 
\] (32)

We deduce that the function \(f_N(\delta, K)\) can now be written as
\[
f_N(\delta, K) = \sum_{n=0}^{N} \sum_{k=0}^{N-n} \left( \frac{1}{2}(p - nq) \right) \left( \frac{1}{K^2} - 1 \right)^k K^{p-nq} f_n \delta^n. 
\] (33)

To first order this expression reduces to
\[
f_1(\delta, K) = \left( 1 - \frac{p}{2} \right) f_0 K^p + \frac{p}{2} f_0 K^{p-2} + f_1 \delta K^{p-q}. 
\] (34)

Applying the principle of least sensitivity \([17]\) leaves us with
\[
\frac{\partial f_1(\delta, K)}{\partial K} \sim p \left( 1 - \frac{p}{2} \right) f_0 + \frac{p(p-2)}{2} f_0 K^{-2} + (p - q) f_1 \delta K^{-q} \equiv 0. 
\] (35)

Next, making the strong-coupling ansatz
\[
K^{(1)}(\delta) = \delta^{1/q} \left( k_0^{(1)} + k_1^{(1)} \delta^{-2/q} + \ldots \right), 
\] (36)
we obtain the following equation from (35):
\[
p \left( 1 - \frac{p}{2} \right) f_0 + \frac{p(p-2)}{2} f_0 (k_0^{(1)})^{1/q}^{-2} + (p - q) f_1 \delta (k_0^{(1)})^{-q} = 0. 
\] (37)

The second term is a subleading contribution in the limit as the coupling \(\delta\) goes to infinity which we can neglect. Solving for \(k_0^{(1)}\) we then get
\[
k_0^{(1)} = \left( \frac{2f_1}{f_0 p(p-2)} \right)^{1/q}. 
\] (38)

Assuming that the ansatz \([\text{30}]\) for the variational parameter \(K(\delta)\) also holds for higher orders we obtain from the function \(f_N(\delta, K)\) in (33)
\[
f_N(\delta) = \delta^{\frac{p}{q}} \left[ b_0^{(N)}(k_0^{(N)}) + b_1^{(N)}(k_0^{(N)}) \delta^{-2/q} + \ldots \right], 
\] (39)
where the leading strong-coupling coefficient \(b_0^{(N)}(k_0^{(N)})\) is given by
\[
b_0^{(N)}(k_0^{(N)}) = \sum_{n=0}^{N-n} \sum_{k=0}^{N-n} \left( \frac{1}{2}(p - nq) \right) (-1)^k f_n (k_0^{(N)})^{p-nq}. 
\] (40)

The inner sum can be further simplified, using
\[
\sum_{k=0}^{m} (-1)^k \binom{\alpha}{k} = (-1)^m \binom{\alpha - 1}{m}. 
\] (41)
Thus the strong-coupling coefficient \( b_0(N) \) reduces to
\[
b_0(N)\left(k_0(N)\right) = \sum_{n=0}^{N}(-1)^{N-n} \left(\frac{1}{2}p - nq \right) f_n\left(k_0(N)\right)^{p-nq}.
\] (42)

So, looking at equation (39) we see that the fraction \( p/q \) tells us the leading power behavior in \( \delta \) and \( 2/q \) indicates the approach to scaling:
\[
\sum_{j=0}^{\infty} f_j \delta^j \sim \delta^{p/q} \left( b_0 + b_1 \delta^{-2/q} + \ldots \right) \quad (\delta \to \infty).
\] (43)

For the instanton equation we can determine the numbers \( p \) and \( q \) by re-obtaining the differential equation (6) from the difference equation (11). The positive real axis is discretized in steps of width \( a \), so that we let \( x_n \equiv na \). The power series expansion for the discrete function \( f_n = f(x_n) \) has the form
\[
f_{n\pm 1} = f(x_n) \pm f'(x_n)a + \frac{1}{2} f''(x_n)a^2 \pm \frac{1}{6} f'''(x_n)a^3 + \frac{1}{24} f''''(x_n)a^4 \pm \ldots.
\] (44)

Thus, the numerator of the second derivative (10) becomes
\[
f_{n+1} - 2f_n + f_{n-1} = f''_n a^2 + \frac{1}{12} f'''_n a^4 + \ldots,
\] (45)

so the zeroth-, first-, and third-order contributions cancel. Translating the lattice result for \( f_n \) back to the continuous function \( f(x_n) = f_n \), the difference equation (11) reads
\[
\epsilon^2 \left[ f''(x) + \frac{1}{12} f''''(x)a^2 + \ldots \right] + f(x) - f^3(x) = 0.
\] (46)

Writing out the power series
\[
f(x) = f_0(x) + a^2 f_1(x) + a^4 f_2(x) + \ldots,
\] (47)
and comparing even powers of \( a \), we get from equation (46) for \( a^0 \n
\epsilon^2 f''_0(x) + f_0(x) - f^3_0(x) = 0,
\] (48)

which is just the original instanton equation (3). For \( a^2 \) we have
\[
\epsilon^2 f''_1(x) + f_1(x) \left( 1 - 3f^2_0(x) \right) = -\frac{1}{12} \epsilon^2 f''''_0(x).
\] (49)

The boundary values read
\[
f_0(0) = 0, \quad f_0(\infty) = 1,
\] (50)
and
\[
f_1(0) = f_1(\infty) = 0,
\] (51)
respectively. The solution to equation (48) with the boundary values (50) is of course
\[
f_0(x) = \tanh \frac{x}{\epsilon \sqrt{2}}
\] (52)
So, finally from (47) we get for the derivative at the origin \( x = 0 \):

\[
f'(0) = f'_0(0) + \frac{\epsilon^2}{\delta} f''_1(0) + \ldots = \frac{1}{\epsilon \sqrt{2}} + \frac{\epsilon^2}{\delta} f''_1(0) + \ldots. 
\]

(53)

Comparing equation (53) with (18), we resum the weak-coupling series in (18) as

\[
1 - \frac{\delta^2}{8} + \ldots = \frac{1}{\sqrt{2}} + \frac{\epsilon^2}{\delta} f''_1(0) \delta + \ldots .
\]

(54)

Also, comparing with (43), we conclude that the leading power and the approach to scaling are given by

\[
\frac{p}{q} = -\frac{1}{2}, \quad \frac{2}{q} = 1,
\]

(55)

respectively. So we identify \( p = -1 \) and \( q = 2 \).

We now evaluate the leading strong-coupling coefficient \( b_0 \) from (43) according to (42) with \( p = -1 \) and \( q = 2 \). To that end we substitute our 200 weak-coupling coefficients from [14] into the formula using a computer algebra program. We are now confronted with the following problem: The principle of least sensitivity cannot be unambiguously applied. Optimizing with respect to extrema, inflection points, or higher derivatives does yield converging results for the strong-coupling limit. However, all these strong-coupling series converge to the wrong values.

There is one particularly unpleasant case: The second derivative with respect to \( k_0 \) for the largest \( k_0 \) where this derivative exists (see Fig. 2) gives a convergent strong-coupling series. The numbers come extremely close to \( 1/\sqrt{2} \) as one can see from the 20 numbers in Table 3. The 200th leading strong-coupling coefficient is \( b^{(200)}_0 = 0.707417 \ldots \). However, a Richardson extrapolation [13] based on the first 200 orders then unfortunately shows that variational perturbation theory produces a value slightly smaller than \( 1/\sqrt{2} \). The first six orders of Richardson extrapolations are presented in Table 4. Hence, the strong-coupling series \( b^{(N)}_0 \) does converge, but it converges to the wrong number, only one part per 1000 away from the true value:

\[
f_1^{(\text{VPT})} \approx \lim_{N \to \infty} \sum_{n=0}^{200} a_{1,n} \delta^n = b^{(\infty)}_0 = 0.7063998320858845 \pm 0.0000000000000001 \]  

(56)

compared with \( f'(0) = 1/\sqrt{2} = 0.7071067812 \ldots \). The deviation is just 0.009\% , but \( 1/\sqrt{2} \) can unfortunately be ruled out.

Given that \( p = -1 \) and \( q = 2 \), the failure of variational perturbation theory is not surprising. According to Ref. [16] the fraction \( 2/q \) must lie within the open interval \( (1/2,1) \). Otherwise, one cannot prove that variational perturbation theory converges. Thus, this problem lies exactly on the upper boundary of the region in which the summation method is known to work.

We can understand the upper edge of the range of the parameter \( 2/q \) that describes the approach to scaling \( 2/q \) by looking at the standard deviation from the actual limiting value. It turns out [16] that the deviation in the limit as the perturbative order \( N \) goes to infinity assumes the shape

\[
\left| \frac{b^{(N)}_0 - b_0}{b_0} \right| \sim \exp \left( -CN^{1-2/q} \right) \quad (N \to \infty),
\]

(57)

where \( C \) is a constant. So, to obtain exponential convergence for the sequence formed by the \( b^{(N)}_0 \), we need \( 1 - 2/q > 0 \). In other words, the approach to scaling \( 2/q \) is bounded and it must be smaller than one. The lower edge is more subtle and is discussed in Ref. [16].
In conclusion, we have applied variational perturbation theory to a case that lies at the very edge of its applicability. We see that variational perturbation theory gives better results by about a factor of 10 than the Padé approximations examined in Sec. II. However, we have not yet found a systematic method for resumming (19) that enables us to perform the continuum limit of the discrete lattice theory. Therefore, we now lay the foundation for further investigations by analyzing the large-order behavior of the instanton series.

IV Large-Order Behavior for the Instanton Equation

It can be seen from the numerical results in [14] that the instanton weak-coupling series is of Borel type. That is, it exhibits an alternating sign pattern. From the ratio test we can see that the coefficients \(a_{n,j}\) do not grow factorially fast. The large-order behavior of \(a_{n,j}\) has the general form

\[
a_{n,j} \sim (-1)^{n+j+1} K_n^j A_n B_n \quad (j \to \infty). \tag{58}
\]

The constant \(A_n\) can be obtained by evaluating the limit

\[
A_n = \lim_{j \to \infty} \frac{\log a_{n,j+2} a_{n,j}}{\log \left(\frac{j(j+2)}{(j+1)^2}\right)}, \tag{59}
\]

and the reciprocal of the radius of convergence is

\[
K_n = - \lim_{j \to \infty} \frac{a_{n,j+1}}{a_{n,j}} \left(\frac{j}{j+1}\right)^{A_n}. \tag{60}
\]

Also, the overall factor \(B_n\) is determined from

\[
B_n = \lim_{j \to \infty} \frac{|a_{n,j}|}{K_n^j A_n}. \tag{61}
\]

Using the 200 weak-coupling coefficients, we find that the exponent \(A_n\) and the reciprocal radius of convergence \(K_n\) are independent of \(n\). The value of \(K_2 = 2.46682906 \) coincides with \(K_1 = 2.46682906\) for all significant digits. The same is true for \(A_1 = -1.500000\) and \(A_2 = -1.500000\). Thus, it appears that we may omit the subscripts \(n\) for \(K_n\) and \(A_n\). In contrast, the data suggests that \(B_n\) strongly depends on \(n\). \(B_n\) is the numerical value associated with the largest uncertainty. In fact, Eq. (61) suggests that small deviations in \(K\) and \(A\) lead to dramatic changes in the value of \(B_n\). We calculated \(A, K, B_1,\) and \(B_2\) up to 200th order with the help of Maple V R7. We then extrapolated these 200 orders to infinity using Richardson extrapolation [13]. We obtained

\[
A = -1.500000 \pm 0.000001, \quad K = 2.46682906 \pm 0.0000001, \quad B_1 = 0.0171 \pm 0.0001, \quad B_2 = 0.1190 \pm 0.0001. \tag{62}
\]

Detailed numerical results for the first six Richardson extrapolations for the exponent \(A\), the inverse radius of convergence \(K\), and the overall factors \(B_1\) and \(B_2\) can be found in Tables [3] and [4]. The calculation of \(B_1\) is extremely delicate; changing the inverse radius of convergence in the sixth decimal place influences the third significant figure of \(B_1\). The same is true of \(B_2\).
Unfortunately, there is no way to derive these values by applying asymptotic analysis to the recursion relation (16). The problem is that the double summation in this equation includes small \( j \), so we cannot let \( j \) go to infinity and use the large-order behavior (58). Substituting the ansatz (58) into equation (16) and taking the limit leads to contradictory results. For \( n = 1 \) we get

\[
K_j A B_1 = \frac{1}{2} (j - 1)^A B_2 + (j - 1)^A B_1 - \frac{3}{2} B_1^2 K \sum_{k=1}^{j-1} k^A (j - k)^A
\]

Pulling out some factors and letting \( x \equiv k/j \), we obtain for the first summation

\[
\lim_{j \to \infty} \sum_{k=1}^{j} \left( \frac{k}{j} \right)^A A^A (1 - \frac{k}{j} )^A = \int_0^1 dx \int_0^1 dy [x(1 - y)]^A = \frac{\Gamma^2(A + 1)}{\Gamma(2A + 2)},
\]

if and only if \( A > -1 \). For \( A < -1 \) which is strongly favored by the data we obtain

\[
\int_0^1 dx \int_0^1 dy [x(1 - y)]^A = 2 \zeta \left( -A \right).
\]

The double summation reduces to

\[
\lim_{j \to \infty} \sum_{k=1}^{j} \sum_{l=1}^{j-k} \frac{k^{A-1}}{j^{2A}} \left( 1 - \frac{k}{j} - \frac{l}{j} \right) = \int_0^1 dx \int_0^1 dy [xy(1 - x - y)]^A = \frac{3 \zeta^2 \left( -A \right)}{\Gamma(3A + 3)},
\]

where \( y \equiv l/j \) and \( A > -1 \). For \( A < -1 \) the result is

\[
\int_0^1 dx \int_0^1 dy [xy(1 - x - y)]^A = 3 \zeta^2 \left( -A \right).
\]

Substituting the results in (65) and (67) into (63) leads to a contradiction: The inverse radius of convergence then turns out to be

\[
K = \frac{1 + \frac{B_1}{2 \pi}}{1 + 3 \zeta \left( \frac{A}{2} \right) B_1 + \frac{3 \zeta^2 \left( \frac{A}{2} \right)}{4} B_1^2},
\]

which would imply that, given \( B_1 = 0.0171 \) and \( B_2 = 0.1190 \), the value of \( K \) would be

\[
K = 3.940.
\]

This result can be ruled out because of the numerical result (62). Also, (68) does not contain the exponent \( A \) because all the factors \( j^A \) in (63) cancel. So \( A \) cannot be determined analytically using this asymptotic analysis.

V Boundary-Layers on the Lattice — Blasius Equation

The Blasius equation [18] arises in the study of fluid dynamics. It is a special limiting case of the Navier-Stokes equation and determines the flow of an incompressible fluid across a semi-infinite flat plate. The equation reads

\[
2xy''(x) + y(x)y''(x) = 0.
\]
Assuming that the tangential velocity \( y'(x) \) at the outer limit of the boundary layer is constant, the boundary conditions read \([19]\)

\[
y(0) = y'(0) = 0, \quad y'(\infty) = 1.
\]  

(71)

Our objective here is to calculate the second derivative \( y''(0) \), which represents the stress on the plate. We discretize the Blasius equation \((70)\) by introducing a lattice spacing \( \delta \):

\[
2\delta(f_{n+1} - 3f_n + 3f_{n-1} - f_{n-2}) + f_n(f_{n+1} - 2f_n + f_{n-1}) = 0,
\]

(72)

where we define \( f_n \equiv y(na)/a \) and \( \delta \equiv \epsilon/a^2 \). The boundary conditions \((71)\) now read

\[
f_0 = f_{-1} = 0, \quad f_n \sim n \quad (n \to \infty).
\]

(73)

Expanding \( f_n \) as a series in powers of \( \delta \) as in Eq. \((14)\), we obtain the recursion relation \([10]\)

\[
a_{n+1,j} - 2a_{n,j} + a_{n-1,j} = -\frac{2}{n} (a_{n+1,j-1} - 3a_{n,j-1} + 3a_{n-1,j-1} - a_{n-2,j-1})
- \frac{1}{n} \sum_{k=1}^{n-1} a_{n,k} (a_{n+1,j-k} - 2a_{n,j-k} + a_{n-1,j-k}),
\]

(74)

The boundary values are

\[
a_{n,0} = n \quad (n \geq 0),
\]

\[
a_{-1,0} = 0,
\]

\[
a_{n,0} = a_{n,j} \quad (n \geq 0).
\]

(75)

Eq. \((74)\) can be solved order by order by using a computer algebra program. Table \([1]\) shows the first 20 weak-coupling coefficients \( a_{1,j} \). All coefficients up to the 300th order can be found at \([20]\).

VI Padé Resummation for the Blasius Equation

We now resum the weak-coupling coefficients using the Padé method \([23]\) with \( M = -1/2 \). This value of \( M \) will be derived in Sec. \( \text{VII} \) in Eq. \((82)\). The exact solution \([10]\) to the Blasius equation \((70)\), obtained numerically up to five digits, is \( y''(0) = 0.33206 \). Unfortunately, the sequence formed by the approximants \( S_N \) appears to converge, but not to the correct value. According to Table \([1]\) the sequence becomes very flat and Richardson extrapolation \([13]\) shows that the \( S_N \) approach the wrong limiting value (see Table \([1]\)). A third-order Richardson gives \( S_\infty = 0.3430 \), based on the first 70 weak-coupling coefficients. This value is significantly higher than the correct value \( y''(0) = 0.33206 \), the deviation is 3.3%.

The failure of the Padé resummation is not surprising because the Padé method assumes the approach to scaling \( \delta^{-1} \) according to \([24]\). However, in the case of the Blasius equation the approach to scaling is \( \delta^{-1/2} \), as we will see in equation \((82)\) in the next section.

VII Variational Perturbation Theory for the Blasius Equation

Variational perturbation theory for the Blasius equation fails to converge to the correct answer in the same way as for the instanton problem. We determined the leading strong-coupling term \([22]\) up to 200th order and again it was impossible to find extrema, inflection
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points, or higher derivatives that yield the correct result. Tables 12 and 13 show the last 20 strong-coupling coefficients $b_k^{(N)}$ and six orders of Richardson extrapolation. By determining the values of $p$ and $q$ we show why variational perturbation is likely to fail for this problem.

Consider again the Taylor expansions for $f_{n\pm 1}$ in (44) together with the Taylor series for $f_{n-2} = f(x_n - 2a)$, namely

$$f_{n-2} = f(x_n) - 2f'(x_n)a + 2f''(x_n)a^2 - \frac{4}{3}f'''(x_n)a^3 + \frac{2}{3}f''''(x_n)a^4 \pm \ldots.$$  (76)

Inserting these expressions into the difference equation for the Blasius problem (72) and translating back to the continuous function $f(x_n) = f_n$, we get

$$2\epsilon \left(f''(x)a - \frac{1}{2}f''''(x)a^2 + \ldots\right) + f(x) \left(f''(x)a^2 + \frac{1}{2}f''''(x)a^4 + \ldots\right) = 0.$$  (77)

Next we transform back to the function $y(x) = af(x)$ and assume the Taylor series

$$y(x) = y_0(x) + ay_1(x) + a^2y_2(x) + \ldots.$$  (78)

To zeroth order in $a$ we obtain

$$2\epsilon y''''(x) + y_0(x)y''(x) = 0,$$  (79)

which is just the Blasius equation (71). The small parameter $a$, which is the lattice spacing, relates $\epsilon$ and $\delta$ by $a = \sqrt{\epsilon/\delta}$. Thus, if we evaluate the Taylor series (78) for the second derivative at the origin, we see that

$$y''(0) = y_0''(0) + ay_1''(0) + \ldots = \frac{0.33206}{\sqrt{\epsilon}} + \frac{\sqrt{\epsilon}}{\delta}y_1''(0) + \ldots.$$  (80)

Comparing this series to the original weak-coupling series

$$y''(0) = \sqrt{\frac{\delta}{\epsilon}} (1 - 2\delta + 2\delta^2 + \ldots),$$  (81)

we can now determine the leading power $p/q$ and the approach to scaling $2/q$:

$$1 - 2\delta + 2\delta^2 + \ldots = \delta^{-1/2} \left(0.33206 + \delta^{-1/2}\epsilon y_1''(0) + \ldots\right),$$  (82)

so we obtain $p = -2$ and $q = 4$.

Again we find that the approach to scaling $2/q = 1/2$ lies just on the boundary of the open interval $(1/2, 1)$, for which the proof of convergence holds. This situation here is the opposite of the instanton case in that it sits at the lower boundary of the open interval in which variational perturbation theory works.

VIII Large-order Behavior for the Blasius Equation

The Blasius equation exhibits a large-order behavior which is a more subtle than for the instanton problem. The Blasius weak-coupling coefficients are not of Borel type; that is, the sign pattern is not alternating. Rather, the sign structure is governed by a cosine function with a frequency that is significantly different from $\pi$. Remarkably, it turns out that a pure cosine $\cos(an)$ cannot reproduce all signs correctly. Up to 300th order the sign structure given by $\cos(an)$ is broken twice: The signs at $n = 62$ and at $n = 212$ are not correct if we optimize with respect to $a$. So we must consider an additional phase shift
cos\((an + b)\). The parameter \(b\) turns out to be slightly smaller than \(\pi\), but it reproduces all 300 signs correctly.

In order to determine the numerical values of \(a\) and \(b\) we define

\[
f(a, b) \equiv \sum_{n=1}^{N} \frac{\cos(an + b)}{\cos(an + b) |a_{1,n}|}.
\]

The sum ends at \(N = 300\) because this is as high as we can calculate using Maple; we know the first 300 weak-coupling coefficients \(a_{1,j}\). For the correct values of \(a\) and \(b\) the function \(f(a, b)\) must be equal to 300. We then plot the function \(f(a, b)\) over the \(a-b\) plane and search for peaks. A careful study of the peaks yields values for \(a\) and \(b\) which allow the function \(f(a, b)\) to assume its maximum at 300. These numbers are given in Table 14.

The large-order behavior of the Blasius weak-coupling coefficients (unlike the large-order behavior of the instanton coefficients) has an additional overall factor \(\cos(an + b)\), and we can now see that the remaining structure differs from the structure of the instanton weak-coupling coefficients. Dividing by the cosine, we observe that the coefficients

\[
a'_{j} \equiv \frac{a_{1,j}}{\cos(a_{j} + b)}
\]

grow factorially fast. Thus, we also divide by \(j!\):

\[
b_{j} \equiv \frac{a_{1,j}}{\cos(a_{j} + b)j!}.
\]

The coefficients \(b_{j}\) are unstable under a ratio test. That is, the ratio \(b_{j+1}/b_{j}\) decreases and then begins to oscillate. This is the inaccuracy that results from the delicate sign pattern of the first 300 coefficients \(a_{1,j}\).

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### Tables

**Table 1:** The first 20 weak-coupling coefficients $a_{1,j}$ for the instanton problem (15) and (16).

| $j$ | $a_{1,j}$ | $j$ | $a_{1,j}$ |
|-----|-----------|-----|-----------|
| 1   | $-\frac{1}{5}$ | 11  | $-2887747$ |
| 2   | $\frac{125}{8}$ | 12  | $93892471$ |
| 3   | 0         | 13  | $-215789295$ |
| 4   | $\frac{11}{128}$ | 14  | $37816709831$ |
| 5   | $-\frac{21}{128}$ | 15  | $-8349041385$ |
| 6   | $\frac{102579}{1024}$ | 16  | $1188129285795$ |
| 7   | $-\frac{5899}{1024}$ | 17  | $-26591064132291$ |
| 8   | $\frac{39203}{32768}$ | 18  | $4789245152559$ |
| 9   | $\frac{8073}{32768}$ | 19  | $108383753179167$ |
| 10  | $\frac{154949}{262144}$ | 20  | $39433620459113981$ |

**Table 2:** The first 20 Padé approximants for the solution to the instanton problem (19).

| $N$ | $S_N$ | $N$ | $S_N$ |
|-----|-------|-----|-------|
| 1   | 0.840896415 | 11  | 0.709998411 |
| 2   | 0.781934407 | 12  | 0.708235422 |
| 3   | 0.757237797 | 13  | 0.706789935 |
| 4   | 0.740759114 | 14  | 0.705659505 |
| 5   | 0.723927185 | 15  | 0.704734605 |
| 6   | 0.731210449 | 16  | 0.704006945 |
| 7   | 0.723927185 | 17  | 0.703419862 |
| 8   | 0.719045188 | 18  | 0.702964717 |
| 9   | 0.715146335 | 19  | 0.702610220 |
| 10  | 0.712308458 | 20  | 0.702349024 |

Table 1: The first 20 weak-coupling coefficients $a_{1,j}$ for the instanton problem (15) and (16).

Table 2: The first 20 Padé approximants for the solution to the instanton problem (19).
### Table 3: The last 20 variational strong-coupling coefficients $b_{0}^{(N)}$ from Eq. (42).

| $N$  | $b_{0}^{(N)}$       |
|------|--------------------|
| 180  | 0.707530492        |
| 181  | 0.707524250        |
| 182  | 0.707518076        |
| 183  | 0.707511970        |
| 184  | 0.707505930        |
| 185  | 0.707499955        |
| 186  | 0.707494044        |
| 187  | 0.707488197        |
| 188  | 0.707482412        |
| 189  | 0.707476687        |
| 190  | 0.707471024        |
| 191  | 0.707465419        |
| 192  | 0.707459872        |
| 193  | 0.707454384        |
| 194  | 0.707448952        |
| 195  | 0.707443575        |
| 196  | 0.707438253        |
| 197  | 0.707432986        |
| 198  | 0.707427771        |
| 199  | 0.707422609        |

### Table 4: Six orders of Richardson extrapolations for the strong-coupling coefficient $b_{0}^{(N)}(k_{0})$ up to $N = 200$ for the instanton problem. The last value is only 0.099% away from the correct limiting value $1/\sqrt{2} = 0.7071067812...$.

| order | value for $b_{0}^{(N)}$       | convergence     |
|-------|--------------------------------|-----------------|
| 1     | 0.70640049                     | decreasing      |
| 2     | 0.70639983200                  | increasing      |
| 3     | 0.706399832082                 | increasing      |
| 4     | 0.7063998320858658             | increasing      |
| 5     | 0.70639983208584411            | increasing      |
| 6     | 0.706399832085846498           | increasing      |

### Table 5: Six orders of Richardson extrapolations for the exponent $A$ of the large-order instanton weak-coupling coefficients, based on the first 200 weak-coupling coefficients. The value $A = -3/2$ is quite plausible.

| order | value for $A$      | convergence     |
|-------|--------------------|-----------------|
| 1     | -1.4998            | increasing      |
| 2     | -1.5000017         | decreasing      |
| 3     | -1.5000011         | decreasing      |
| 4     | -1.49999874        | increasing      |
| 5     | -1.5000004         | decreasing      |
| 6     | -1.499999893       | increasing      |
Table 6: Six orders of Richardson extrapolations for the inverse radius of convergence $K$ of the large-order instanton weak-coupling coefficients, based on the first 200 weak-coupling coefficients under the assumption that $A = -3/2$.

| order | value for $K$ | convergence |
|-------|---------------|-------------|
| 1     | 2.46692       | decreasing  |
| 2     | 2.4668283     | increasing  |
| 3     | 2.46682911    | decreasing  |
| 4     | 2.466829065   | decreasing  |
| 5     | 2.4668290597  | increasing  |
| 6     | 2.4668290635  | decreasing  |

Table 7: Six orders of Richardson extrapolations for the overall factor $B_1$ of the large-order instanton weak-coupling coefficients, based on the first 200 weak-coupling coefficients under the assumption that $K = 2.4482906$ and $A = -3/2$. The value of $B_1$ strongly depends on the numerical values for $A$ and $K$. Changing $K$ in the sixth decimal place influences the third significant figure of $B_1$. Also, all the Richardson extrapolations are increasing so, strictly speaking, we only have a lower boundary for $B_1$. Thus, the accuracy of $B_1$ may not be very good.

| order | value for $B_1$ | convergence |
|-------|----------------|-------------|
| 1     | 0.0170837      | increasing  |
| 2     | 0.0170864      | increasing  |
| 3     | 0.017087       | increasing  |
| 4     | 0.0170893      | increasing  |
| 5     | 0.0170908      | increasing  |
| 6     | 0.0170922      | increasing  |

Table 8: Six orders of Richardson extrapolations for the overall factor $B_2$ of the large-order instanton weak-coupling coefficients based on the first 200 weak-coupling coefficients and the same assumptions as in the case of $B_1$ (see Table 7). The value of $B_2$ depends strongly on $A$ and $K$.

| order | value for $B_2$ | convergence |
|-------|----------------|-------------|
| 1     | 0.119069       | increasing  |
| 2     | 0.119083       | increasing  |
| 3     | 0.119093       | increasing  |
| 4     | 0.119054095    | increasing  |
| 5     | 0.119054125    | increasing  |
| 6     | 0.119054146    | increasing  |
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Table 9: The first 20 weak-coupling coefficients for the Blasius recursion relation (74) and (75). Observe that the coefficients \( a_{1,j} \) are not of Borel type (they do not alternate in sign). A cosine function with a frequency different from \( \pi \) governs the sign pattern (see Sec. VIII).

| \( j \) | \( a_{1,j} \) | \( j \) | \( a_{1,j} \) |
|---|---|---|---|
| 1 | -2 | 11 | 30868632383 \( \frac{5437375}{5437375} \) |
| 2 | 2 | 12 | 6325029622 \( \frac{637875}{637875} \) |
| 3 | \( \frac{8}{9} \) | 13 | -48769375019181 \( \frac{13408770375}{13408770375} \) |
| 4 | -6 | 14 | -4774319527974167 \( \frac{37819608750}{37819608750} \) |
| 5 | -184/15 | 15 | 43032125108745734 \( \frac{2212447111875}{2212447111875} \) |
| 6 | \( \frac{136}{9} \) | 16 | 796235444548876905317 \( \frac{6399861541875}{6399861541875} \) |
| 7 | 11062/105 | 17 | -2249988054508764174584049 \( \frac{6776855230149837500}{6776855230149837500} \) |
| 8 | \( \frac{838}{235} \) | 18 | -178060575619150189817796 \( \frac{14237097164915625}{14237097164915625} \) |
| 9 | -15557451/14175 | 19 | 13224896152219729367496038639 \( \frac{1301999758346024437500}{1301999758346024437500} \) |
| 10 | \( \frac{57629622}{99225} \) | 20 | 121756993154067534451732128947029 \( \frac{1153217988475534737500}{1153217988475534737500} \) |

Table 10: The first 20 Padé approximants for the solution to the Blasius equation (70). The sequence formed by the \( S_N \) converges extremely slowly.

| \( N \) | \( S_N \) | \( N \) | \( S_N \) |
|---|---|---|---|
| 1 | 0.3445 | 11 | 0.3574632121 |
| 2 | 0.4204482076 | 12 | 0.3563326651 |
| 3 | 0.3982018303 | 13 | 0.3553848048 |
| 4 | 0.3819443732 | 14 | 0.354579944 |
| 5 | 0.3742062309 | 15 | 0.3538882842 |
| 6 | 0.3690504811 | 16 | 0.3532891509 |
| 7 | 0.3653779673 | 17 | 0.3527655813 |
| 8 | 0.3626359060 | 18 | 0.3523046588 |
| 9 | 0.3605155915 | 19 | 0.3518961929 |
| 10 | 0.3588309707 | 20 | 0.3515320399 |

Table 11: Three orders of Richardson extrapolations for the Blasius equation (70), based on the first 70 Padé approximants \( S_N \).

| order | value of \( y''(0) \) | convergence |
|---|---|---|
| 1 | 0.3445 | decreasing |
| 2 | 0.3436 | decreasing |
| 3 | 0.3430 | oscillating |
Table 12: The last 20 variational strong-coupling coefficients $b_0^{(N)}$ for the Blasius equation. The very last coefficient is $b_0^{(200)} = 0.33695931237713$, as opposed to the correct value $y''(0) = 0.33206$.

| $N$ | $b_0^{(N)}$ | $N$ | $b_0^{(N)}$ |
|-----|-------------|-----|-------------|
| 180 | 0.33696017793094 | 190 | 0.33695971119646 |
| 181 | 0.33696012777085 | 191 | 0.33695966849139 |
| 182 | 0.33696007843082 | 192 | 0.33695962644843 |
| 183 | 0.33696002989308 | 193 | 0.33695958505396 |
| 184 | 0.33695998214034 | 194 | 0.33695954429471 |
| 185 | 0.33695993515575 | 195 | 0.33695950415774 |
| 186 | 0.33695988892292 | 196 | 0.33695946463046 |
| 187 | 0.33695984342591 | 197 | 0.33695942570058 |
| 188 | 0.33695979864918 | 198 | 0.33695938735612 |
| 189 | 0.33695975457760 | 199 | 0.33695934958540 |

Table 13: Six orders of Richardson extrapolations for the strong-coupling coefficient $b_0^{(N)}(k_0)$ up to $N = 200$ for the Blasius equation. The last value is 1.5% away from the correct limiting value $y''(0) = 0.33206$.

| order | value for $b_0^{(N)}$ | convergence |
|-------|------------------------|-------------|
| 1     | 0.3369518              | increasing  |
| 2     | 0.336955563            | increasing  |
| 3     | 0.336955600539         | increasing  |
| 4     | 0.3369556008803        | increasing  |
| 5     | 0.336955600883462      | increasing  |
| 6     | 0.33695560088349232    | increasing  |

Table 14: Examples of the parameters $a$ and $b$ that give the first 300 signs of the Blasius weak-coupling coefficients correctly, assuming that the sign structure of the underlying large-order behavior is of the form $\cos(an + b)$. The last two values for $a$ can be obtained approximately by summing $2\pi$ to the first two values.

| $a$   | $b$   |
|-------|-------|
| 1.3941| 3.09  |
| 1.3939| 3.11  |
| 7.67830| 3.031 |
| 7.67686| 3.130 |
Figures

Figure 1: The real part of the Padé approximants $S_N$ up to 200th order. Note that the approximants do not converge to the exact solution, which is represented by the horizontal solid line. The phases where the approximants become complex are marked by spikes.
Figure 2: The function $b_0^{(N)}(k_0^{(N)})$ from (42) for $N = 200$ (solid line) and its second derivative with respect to $k_0$ (dotted line). The upper horizontal line equals $1/\sqrt{2}$, the correct limiting value of the instanton problem. All extrema of $b_0^{(N)}$ are far from this value. Only the inflection point on the right-hand side comes close. The value for $k_0$, for which the second derivative vanishes, is $k_0 = 18.42510$. Substituting that number into the function $b_0^{(N)}(k_0^{(N)})$, we obtain in the 200th order $b_0^{(200)} = 0.707417$. The corresponding Richardson extrapolations can be found in Table 4.