

Article

Dynamic Response of a Tunnel Lining Under the Coupled Influence of Near-Field Effects and an Elastic–Slip Interface

Yao Rong ^{1,2}, Junping Yu ^{1,2,3}, Zhiyun Liu ^{3,*}, Jingliang Dong ³ and Yongwei Li ³¹ Jiangxi Communications Investment Maintenance Technology Group Co., Ltd., Nanchang 330200, China² Jiangxi Provincial Key Laboratory of Highway Bridge and Tunnel Engineering, Nanchang 330200, China³ State Key Laboratory of Safety and Resilience of Civil Engineering in Mountain Area, East China Jiaotong University, Nanchang 330013, China

* Correspondence: zyl15172@126.com

Abstract

This paper presents an analytical solution for near-field cylindrical P-wave scattering by a circular tunnel. A comprehensive elastic–slip interface model is incorporated to capture localized discontinuous deformations and frictional sliding. Employing the wave function expansion method and Graf’s addition theorem, the dynamic stress concentration factor (DSCF) of both soft and stiff linings is systematically investigated. The results reveal that near-field wavefront curvature severely amplifies stress concentration on the illuminated side, degenerating into a plane-wave response only when the dimensionless source distance exceeds 50. Furthermore, within the parameter range considered, the imperfect interface exhibits a dual “flexible barrier” effect heavily influenced by impedance matching: it tends to act as a seismic isolator that reduces the DSCF for stiff linings, but can weaken boundary confinement and exacerbate stress concentration for soft linings. These findings provide crucial theoretical insights for the blast-resistant design of underground structures.

Keywords: cylindrical P-waves; Near-Field Effects; elastic–slip interface; tunnel lining; wave scattering

1. Introduction

Underground infrastructure, such as transportation tunnels [1] and mining roadways [2], plays an indispensable supporting role in modern geotechnical engineering. Despite the conventional consensus that subterranean structures are inherently secure due to the continuous confinement provided by surrounding geological strata, severe dynamic disturbances induced by adjacent drilling-and-blasting operations or seismic events frequently precipitate lining cracking [3], spalling, and even catastrophic structural failure [4]. Consequently, accurately evaluating the dynamic response of lined tunnels under stress waves—particularly the dynamic stress concentration factor (DSCF)—is essential for ensuring long-term structural safety.

In classical elastodynamic derivations, the physical boundary between the hosting rock mass and the tunnel lining is traditionally idealized as a perfectly bonded interface [5–8]. However, constrained by practical construction techniques and the inevitable existence of joint fissures, micro-cracks, or thin soft interlayers, the genuine contact zone universally exhibits distinct imperfect mechanical behaviors [9]. Previous researchers have extensively applied linear spring models to characterize the discrete displacement discon-

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tinuities (jumps) across these imperfect boundaries [10–12]. For instance, Yi et al. [13] investigated the influence of imperfect interfaces on tunnel linings subjected to plane P-waves, demonstrating that boundary imperfection significantly exacerbates the localized DSCF. Furthermore, Zhang et al. [14] explored the impact of multi-layered imperfect interfaces on composite linings and revealed a distinct compounding amplification effect along with novel stress redistribution features. Although the aforementioned spring-type models successfully capture the elastic restoration capacity of the contact zone, they inherently overlook the irreversible tangential friction sliding and interfacial shear degradation triggered by intense dynamic loads [15–17]. To bridge this theoretical gap, advanced elastic-slip interface models have been introduced. By mathematically coupling the interfacial spring stiffness with a localized shear slip coefficient, this advanced framework provides a rigorous continuum-mechanical representation of genuine interfacial mismatch and frictional energy dissipation during dynamic interactions [7].

Beyond the inherent complexities of boundary constitutive behavior, the spatial geometric topography of the incident wavefront fundamentally dictates the scattering and diffraction mechanisms around underground cavities. A vast majority of classical analytical solutions rely strictly on the plane-wave assumption [12,18], which retains physical validity only when the dynamic disturbance source is located at an infinite distance (i.e., the far-field equivalent). In realistic engineering blasting or close-proximity excavations, however, the detonation epicenter is invariably situated close to the existing tunnel structure [6]. Under such near-field configurations, the induced dynamic stress field manifests as cylindrical waves with a pronounced wavefront curvature and pronounced geometric spreading attenuation [19]. Neglecting these near-field attributes and oversimplifying the incident field into plane waves will inevitably underestimate the severe localized energy focusing on the blast-facing surface, thereby introducing significant hidden hazards into the safety evaluation of the lining structure [20,21]. While recent advancements have successfully explored deploying surface resonators to mitigate seismic wave propagation and protect structures [22–24], the dynamic response of deep underground tunnel linings subjected to nearby blasting presents a distinctly different and highly localized challenge.

Recognizing these coupled theoretical gaps, the central question addressed by this research is the analytical determination of how the wavefront curvature of near-field cylindrical compressional waves interacts in a coupled manner with a mechanically imperfect rock–lining interface. The primary objective is to evaluate how this coupling governs stress redistribution and alters the DSCF in the tunnel lining, explicitly distinguishing the dynamic effects between “soft” and “stiff” support structures [25–27]. To achieve this, an analytical framework is established. While relying on classical elastodynamic methodologies—specifically the wave function expansion method combined with Graf’s addition theorem [28–30]—the theoretical contribution of this work lies in the rigorous mathematical coupling of near-field geometric spreading effects with discontinuous elastic-slip boundary conditions. This approach effectively decouples the wave scattering and transmission mechanisms across the dual-layer media. Subsequently, a systematic parametric study is performed to clarify the coupled regulatory mechanisms of the blast source distance, interface compliance, and tangential slip behaviors on the structural DSCF. By comprehensively comparing the distinct wave interference patterns, this research aims to provide a solid and exact analytical benchmark for the blast-resistant design of deep underground engineering.

2. Mechanical Model Considering Slip Characteristics and Incident Cylindrical Waves

2.1. Physical and Geometric Model

Referring to the schematic diagram illustrated in Figure 1, the spatial domain of the tunnel–soil system is mapped using dual coordinate systems to precisely capture the spatial relationship between the blast source and the underground structure. A global Cartesian coordinate system (x, y) and its corresponding polar coordinate system (θ, r) are established at the geometric center of the circular tunnel. Herein, the geometric boundaries of the tunnel lining are defined by the inner radius a and outer radius b . To incorporate the localized effects of the detonation, a secondary, localized polar coordinate system (θ', r') is introduced at the blast source epicenter, which is positioned at a standoff distance L from the tunnel center. Given that the incident waves propagate orthogonally to the longitudinal axis of the tunnel, the dynamic interaction problem is reasonably simplified to a two-dimensional (2D) plane strain state.

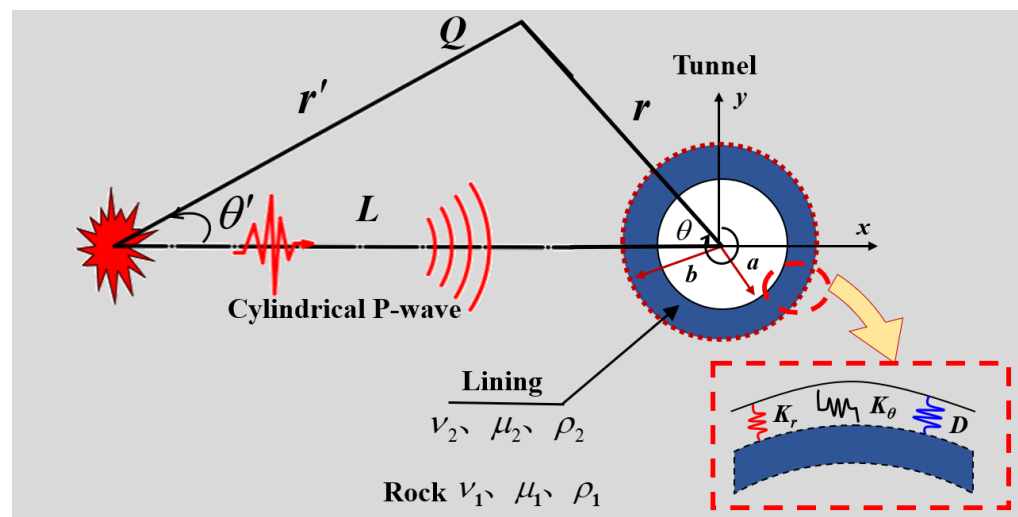


Figure 1. Physical and geometric model.

2.2. Wavefield Representations in the Soil and Lining

In this study, both the surrounding soil medium and the tunnel lining are modeled as single-phase, homogeneous, isotropic, and linear elastic media. In the localized polar coordinate system (θ', r') centered at the blast epicenter, the incident cylindrical P-wave generated by the detonation source can be mathematically expressed as:

$$\varphi_i = \varphi_0 H_0^{(1)}(\alpha_1 r') e^{-i\alpha x} \quad (1)$$

where φ_0 represents the initial amplitude of the incident wave potential, and α_1 is the compressional wavenumber of the surrounding rock matrix. According to the law of cosines, the actual geometric distance from the wave source to the calculation point is given by $r' = \sqrt{r^2 + L^2 + 2rL \cos \theta}$. The function $H_0^{(1)}(\cdot)$ inherently satisfies the Sommerfeld radiation condition for outward-propagating cylindrical waves. For simplicity in the steady-state analysis, the time-harmonic factor $e^{-i\omega t}$ is uniformly omitted in all subsequent derivations.

To strictly apply the boundary conditions at the tunnel–rock interface ($r = b$), it is mathematically imperative to transfer the incident wave potential into the global polar coordinate system (θ, r) centered at the tunnel. By employing Graf's addition theorem

and considering the dynamic response domain $r < L$, the incident wave potential φ_i can be expanded into a series:

$$\varphi_i = \varphi_0 \sum_{n=0}^{\infty} \varepsilon_n (-1)^n H_n^{(1)}(\alpha_1 L) J_n(\alpha_1 r) \cos(n\theta) \quad (2)$$

where ε_n is the Neumann factor ($\varepsilon_0 = 1$; $\varepsilon_n = 2$ for $n \geq 1$), and $J_n(\cdot)$ is the Bessel function of the first kind.

Upon striking the tunnel structure, the incident wave undergoes scattering, generating outward-propagating reflected P-waves and SV-waves in the surrounding rock. Therefore, the total wave potential fields in the surrounding rock matrix are formulated as the superposition of the incident and scattered fields:

$$\begin{aligned} \varphi^{(1)} &= \sum_{n=0}^{\infty} \left[\varphi_0 \varepsilon_n (-1)^n H_n^{(1)}(\alpha_1 L) J_n(\alpha_1 r) + A_n H_n^{(1)}(\alpha_1 r) \right] \cos(n\theta) \\ \psi^{(1)} &= \sum_{n=0}^{\infty} B_n H_n^{(1)}(\beta_1 r) \sin(n\theta) \end{aligned} \quad (3)$$

where β_1 is the shear wavenumber of the surrounding rock, and A_n and B_n denote the unknown scattering expansion coefficients.

Within the finite domain of the single-layer tunnel lining, wave propagation exhibits complex multiple reflections between the outer interface ($r = b$) and the inner free surface ($r = a$). Consequently, the comprehensive potential fields for the lining must incorporate both outward-propagating ($H_n^{(1)}(\cdot)$) and inward-propagating ($H_n^{(2)}(\cdot)$) wave components:

$$\begin{aligned} \varphi^{(2)} &= \sum_{n=0}^{\infty} \left[C_n H_n^{(1)}(\alpha_2 r) + D_n H_n^{(2)}(\alpha_2 r) \right] \cos(n\theta) \\ \psi^{(2)} &= \sum_{n=0}^{\infty} \left[E_n H_n^{(1)}(\beta_2 r) + F_n H_n^{(2)}(\beta_2 r) \right] \sin(n\theta) \end{aligned} \quad (4)$$

where α_2 and β_2 are the compressional and shear wavenumbers of the lining, respectively. The terms C_n, D_n, E_n , and F_n constitute the set of undetermined transmission coefficients.

2.3. Formulations of Elastic–Slip Interface and Boundary Conditions

To accurately characterize the mechanical behavior at the contact non-uniformity between the surrounding rock mass and the tunnel lining, an advanced elastic–slip interface model is introduced herein. In this framework, the stress fields across the interface remain strictly continuous, whereas the displacement fields exhibit discrete discontinuities (jumps) governed by the interface flexibility. At the outer lining boundary ($r = b$), the normal and shear stress continuity conditions are dynamically expressed as:

$$\sigma_{rr}^{(1)} = \sigma_{rr}^{(2)}, \quad \tau_{r\theta}^{(1)} = \tau_{r\theta}^{(2)} \quad (5)$$

where the superscripts (1) and (2) designate the surrounding rock matrix (Medium 1) and the single-layer tunnel lining (Medium 2), respectively.

Concurrently, the radial displacement jump across the imperfect interface is modeled via linear spring stiffness, which is proportional to the local normal stress:

$$u_r^{(1)} - u_r^{(2)} = \frac{1}{K_r} \sigma_{rr}^{(2)} \quad (6)$$

where K_r represents the dimensionless radial interface spring stiffness.

More importantly, to incorporate the tangential friction sliding and interfacial shear degradation induced by nearby blasting disturbances, a slip coefficient D is coupled

with the tangential spring stiffness K_θ . Consequently, the modified tangential displacement jump condition at $r = b$ is defined as:

$$u_\theta^{(1)} - u_\theta^{(2)} = \left(\frac{1}{K_\theta} + D \right) \tau_{r\theta}^{(2)} = \bar{S}_t \tau_{r\theta}^{(2)} \quad (7)$$

where $\bar{S}_t = 1/K_\theta + D$ denotes the comprehensive tangential flexibility coefficient of the interface, seamlessly capturing both elastic deformation and frictional energy dissipation mechanisms.

Furthermore, the inner surface of the tunnel lining ($r = a$) is assumed to be an ideal traction-free boundary, yielding the following equations:

$$\sigma_{rr}^{(2)} = 0, \quad \tau_{r\theta}^{(2)} = 0 \quad (8)$$

By substituting the wave potential series into the aforementioned boundary conditions at $r = b$ and $r = a$, a comprehensive 6×6 algebraic matrix equation can be successfully established to determine the unknown wave coefficients.

2.4. Resolution of Matrix Equation

To solve for the coefficients, we first present the stress and displacement equations, and then substitute the boundary conditions to obtain the coefficients. According to the elastodynamic theory, the macroscopic mechanical response of the medium—specifically, the displacement vectors and the Cauchy stress tensors in the polar coordinate system (θ, r) —can be rigorously derived from the scalar potential φ and the vector potential ψ . The geometric gradient relations and linear elastic constitutive laws are given by:

$$\begin{cases} u_r = \frac{\partial \varphi}{\partial r} + \frac{1}{r} \frac{\partial \psi}{\partial \theta} \\ u_\theta = \frac{1}{r} \frac{\partial \varphi}{\partial \theta} - \frac{\partial \psi}{\partial r} \\ \sigma_{rr} = \lambda \nabla^2 \varphi + 2\mu \left[\frac{\partial^2 \varphi}{\partial r^2} + \frac{\partial}{\partial r} \left(\frac{1}{r} \frac{\partial \psi}{\partial \theta} \right) \right] \\ \sigma_{\theta\theta} = \lambda \nabla^2 \varphi + 2\mu \left[\frac{1}{r} \left(\frac{\partial \varphi}{\partial r} + \frac{1}{r} \frac{\partial^2 \varphi}{\partial \theta^2} \right) + \frac{1}{r} \left(\frac{1}{r} \frac{\partial \psi}{\partial \theta} - \frac{\partial}{\partial r} \frac{\partial \psi}{\partial \theta} \right) \right] \\ \sigma_{r\theta} = 2\mu \left(\frac{1}{r} \frac{\partial^2 \varphi}{\partial r \partial \theta} - \frac{1}{r^2} \frac{\partial \varphi}{\partial \theta} \right) + \mu \left[\frac{1}{r^2} \frac{\partial^2 \psi}{\partial \theta^2} - r \frac{\partial}{\partial r} \left(\frac{1}{r} \frac{\partial \psi}{\partial r} \right) \right] \end{cases} \quad (9)$$

where λ and μ denote the Lamé constants of the respective elastic medium.

By substituting the expanded Fourier–Bessel wave potential series into the aforementioned differential operators, the steady-state dynamic fields can be obtained.

For the surrounding rock mass (Medium 1), incorporating the cylindrical geometric spreading factor of the incident wave, the corresponding stress components are expanded as follows:

$$\begin{aligned} \sigma_{rr}^{(1)} &= \frac{2\mu_1}{r^2} \sum_{n=0}^{\infty} \left[I_n \varepsilon_{11}^{(1)}(\alpha_1 r) + A_n \varepsilon_{11}^{(3)}(\alpha_1 r) + B_n \varepsilon_{12}^{(3)}(\beta_1 r) \right] \cos(n\theta) \\ \sigma_{\theta\theta}^{(1)} &= \frac{2\mu_1}{r^2} \sum_{n=0}^{\infty} \left[I_n \varepsilon_{21}^{(1)}(\alpha_1 r) + A_n \varepsilon_{21}^{(3)}(\alpha_1 r) + B_n \varepsilon_{22}^{(3)}(\beta_1 r) \right] \cos(n\theta) \\ \tau_{r\theta}^{(1)} &= \frac{2\mu_1}{r^2} \sum_{n=0}^{\infty} \left[I_n \varepsilon_{41}^{(1)}(\alpha_1 r) + A_n \varepsilon_{41}^{(3)}(\alpha_1 r) + B_n \varepsilon_{42}^{(3)}(\beta_1 r) \right] \sin(n\theta) \end{aligned} \quad (10)$$

where $I_n = \varphi_0 \varepsilon_n (-1)^n H_n^{(1)}(\alpha_1 L)$ represents the modulation coefficient of the near-field incident cylindrical P-wave.

Similarly, within the finite domain of the single-layer tunnel lining (Medium 2), the dynamic stresses and displacements are formulated by superimposing the inward- and outward-propagating wave components:

$$\begin{aligned}\sigma_{rr}^{(2)} &= \frac{2\mu_2}{r^2} \sum_{n=0}^{\infty} [C_n \varepsilon_{11}^{(3)}(\alpha_2 r) + D_n \varepsilon_{11}^{(4)}(\alpha_2 r) + E_n \varepsilon_{12}^{(3)}(\beta_2 r) + F_n \varepsilon_{12}^{(4)}(\beta_2 r)] \cos(n\theta) \\ \sigma_{\theta\theta}^{(2)} &= \frac{2\mu_2}{r^2} \sum_{n=0}^{\infty} [C_n \varepsilon_{21}^{(3)}(\alpha_2 r) + D_n \varepsilon_{21}^{(4)}(\alpha_2 r) + E_n \varepsilon_{22}^{(3)}(\beta_2 r) + F_n \varepsilon_{22}^{(4)}(\beta_2 r)] \cos(n\theta) \quad (11)\end{aligned}$$

$$\begin{aligned}\tau_{r\theta}^{(2)} &= \frac{2\mu_2}{r^2} \sum_{n=0}^{\infty} [C_n \varepsilon_{41}^{(3)}(\alpha_2 r) + D_n \varepsilon_{41}^{(4)}(\alpha_2 r) + E_n \varepsilon_{42}^{(3)}(\beta_2 r) + F_n \varepsilon_{42}^{(4)}(\beta_2 r)] \sin(n\theta) \\ u_r^{(2)} &= \frac{1}{r} \sum_{n=0}^{\infty} [C_n \varepsilon_{71}^{(3)}(\alpha_2 r) + D_n \varepsilon_{71}^{(4)}(\alpha_2 r) + E_n \varepsilon_{72}^{(3)}(\beta_2 r) + F_n \varepsilon_{72}^{(4)}(\beta_2 r)] \cos(n\theta) \\ u_\theta^{(2)} &= \frac{1}{r} \sum_{n=0}^{\infty} [C_n \varepsilon_{81}^{(3)}(\alpha_2 r) + D_n \varepsilon_{81}^{(4)}(\alpha_2 r) + E_n \varepsilon_{82}^{(3)}(\beta_2 r) + F_n \varepsilon_{82}^{(4)}(\beta_2 r)] \sin(n\theta) \quad (12)\end{aligned}$$

The determination of the unknown scattering and transmission coefficients heavily relies on the precise enforcement of the boundary conditions. By invoking the orthogonal characteristics of the Fourier series with respect to the polar angle θ , the wave potential equations are transformed into a set of independent linear algebraic equations for each truncation order n .

The comprehensive governing mechanics of the tunnel–soil system are systematically formulated into a unified matrix representation. By considering the coefficients of the cylindrical wave and the elastic–slip interface, all analytical elements are explicitly populated within the matrix structure:

$$\begin{bmatrix} \varepsilon_{11}^{(3)}(\alpha_1 b) & \varepsilon_{12}^{(3)}(\beta_1 b) & -\frac{1}{\bar{\mu}} \varepsilon_{11}^{(3)}(\alpha_2 b) & -\frac{1}{\bar{\mu}} \varepsilon_{11}^{(4)}(\alpha_2 b) & -\frac{1}{\bar{\mu}} \varepsilon_{12}^{(3)}(\beta_2 b) & -\frac{1}{\bar{\mu}} \varepsilon_{12}^{(4)}(\beta_2 b) \\ \varepsilon_{41}^{(3)}(\alpha_1 b) & \varepsilon_{42}^{(3)}(\beta_1 b) & -\frac{1}{\bar{\mu}} \varepsilon_{41}^{(3)}(\alpha_2 b) & -\frac{1}{\bar{\mu}} \varepsilon_{41}^{(4)}(\alpha_2 b) & -\frac{1}{\bar{\mu}} \varepsilon_{42}^{(3)}(\beta_2 b) & -\frac{1}{\bar{\mu}} \varepsilon_{42}^{(4)}(\beta_2 b) \\ \varepsilon_{71}^{(3)}(\alpha_1 b) & \varepsilon_{72}^{(3)}(\beta_1 b) & M_{33} & M_{34} & M_{35} & M_{36} \\ \varepsilon_{81}^{(3)}(\alpha_1 b) & \varepsilon_{82}^{(3)}(\beta_1 b) & M_{43} & M_{44} & M_{45} & M_{46} \\ 0 & 0 & \varepsilon_{11}^{(3)}(\alpha_2 a) & \varepsilon_{11}^{(4)}(\alpha_2 a) & \varepsilon_{12}^{(3)}(\beta_2 a) & \varepsilon_{12}^{(4)}(\beta_2 a) \\ 0 & 0 & \varepsilon_{41}^{(3)}(\alpha_2 a) & \varepsilon_{41}^{(4)}(\alpha_2 a) & \varepsilon_{42}^{(3)}(\beta_2 a) & \varepsilon_{42}^{(4)}(\beta_2 a) \end{bmatrix} \begin{Bmatrix} A_n \\ B_n \\ C_n \\ D_n \\ E_n \\ F_n \end{Bmatrix} = \begin{Bmatrix} Y_1 \\ Y_2 \\ Y_3 \\ Y_4 \\ 0 \\ 0 \end{Bmatrix} \quad (13)$$

where $\bar{\mu} = \mu_1 / \mu_2$ is the ratio of the shear moduli of the rock mass and lining. The load vector elements $\{Y_i\}$ on the right-hand side, which are induced by the incident near-field cylindrical P-wave, are explicitly derived as $Y_1 = -\varepsilon_n(-1)^n H_n^{(1)}(\alpha_1 L) \varepsilon_{11}^{(1)}(\alpha_1 b)$, $Y_2 = -\varepsilon_n(-1)^n H_n^{(1)}(\alpha_1 L) \varepsilon_{41}^{(1)}(\alpha_1 b)$, $Y_3 = -\varepsilon_n(-1)^n H_n^{(1)}(\alpha_1 L) \varepsilon_{71}^{(1)}(\alpha_1 b)$, and $Y_4 = -\varepsilon_n(-1)^n H_n^{(1)}(\alpha_1 L) \varepsilon_{81}^{(1)}(\alpha_1 b)$. Additionally, to rigorously capture the discontinuous displacement across the imperfect elastic–slip interface, the highly coupled matrix elements $M_{33} \sim M_{36}$ (representing the radial spring stiffness conditions) and $M_{43} \sim M_{46}$ (representing the tangential slip conditions) are fully expanded as follows: $M_{33} = -[\varepsilon_{71}^{(3)}(\alpha_2 b) + 1/K_r \cdot \varepsilon_{11}^{(3)}(\alpha_2 b)]$, $M_{34} = -[\varepsilon_{71}^{(4)}(\alpha_2 b) + 1/K_r \cdot \varepsilon_{11}^{(4)}(\alpha_2 b)]$, $M_{35} = -[\varepsilon_{72}^{(3)}(\beta_2 b) + 1/K_r \cdot \varepsilon_{12}^{(3)}(\beta_2 b)]$, $M_{36} = -[\varepsilon_{72}^{(4)}(\beta_2 b) + 1/K_r \cdot \varepsilon_{12}^{(4)}(\beta_2 b)]$, $M_{43} = -[\varepsilon_{81}^{(3)}(\alpha_2 b) + \bar{S}_t \varepsilon_{41}^{(3)}(\alpha_2 b)]$, $M_{44} = -[\varepsilon_{81}^{(4)}(\alpha_2 b) + \bar{S}_t \varepsilon_{41}^{(4)}(\alpha_2 b)]$, $M_{45} = -[\varepsilon_{82}^{(3)}(\beta_2 b) + \bar{S}_t \varepsilon_{42}^{(3)}(\beta_2 b)]$, and $M_{46} = -[\varepsilon_{82}^{(4)}(\beta_2 b) + \bar{S}_t \varepsilon_{42}^{(4)}(\beta_2 b)]$. To concisely represent these sophisticated mathematical series, a comprehensive set of dimensionless wave operators $\varepsilon_{ij}^{(l)}(\cdot)$ is introduced. The unabridged analytical formulations for these basic operators are provided in Appendix A.

Holes induce stress concentration and deformation in the surrounding rock and lining under stress wave impact. In this paper, the dynamic stress concentration factor is defined as the absolute value of the ratio of circumferential stress with holes to the initial circumferential stress of the intact structure, which is expressed as $DSCF = |\sigma_{\theta\theta} / \sigma_0|$. It is important to note that in practical blasting engineering, detonations generate highly transient shock waves characterized by short-duration pulses and rapid peak pressure decay. The steady-state harmonic formulation adopted in this study essentially provides the fundamental frequency-domain transfer function of the rock–lining system. Based on linear system theory, the transient dynamic response under realistic blast loading can be systematically synthesized by decomposing the incident transient pulse via the Fourier transform, computing the steady-state response across the frequency spectrum, and subsequently applying an inverse Fourier superposition. Therefore, the steady-state solutions derived herein serve as the indispensable analytical foundation for evaluating complex transient blast scenarios.

3. Model Validation and Degeneration

3.1. Non-Dimensionalization Framework

To eliminate the dimensional constraints imposed by specific geometric sizes and physical units, and to significantly enhance the universal applicability of the derived analytical solution across various engineering scales, a systematic non-dimensionalization modeling framework is introduced. By setting the tunnel inner radius a and the shear modulus of the surrounding rock matrix μ_1 as the fundamental scaling benchmarks, the core physical variables are transformed into dimensionless parameters. Accordingly, the dimensionless compressional wavenumber is defined as $\alpha_1 a = \omega a / C_{p1}$, which characterizes the coupling relationship between the dynamic excitation frequency and the structural scale. The relative spatial proximity of the explosive source is quantified by the dimensionless distance L/a . Furthermore, to construct a generalized boundary governing formulation, the mechanical properties of the interface are mapped into the dimensionless normal spring stiffness $K_r^* = K_r / (\mu_1 / a)$, the dimensionless tangential spring stiffness $K_\theta^* = K_\theta / (\mu_1 / a)$, and the dimensionless tangential slip coefficient $D^* = D \cdot (\mu_1 / a)$, respectively. This framework inherently ensures scale consistency during the parametric sensitivity analysis. This study systematically investigates two distinct lining scenarios: the soft lining configuration and the stiff lining configuration. The specific physical and mechanical parameters for both the surrounding rock mass and the tunnel lining are detailed in Table 1. This precise agreement thoroughly substantiates the exactitude of the theoretical analytical solution derived in this study, confirming its reliability for subsequent theoretical investigations.

Table 1. The values of dimensionless parameters for different media [21].

Medium	Case 1	Case 2
C_{p1}/C_{p2}	1.5	0.7
μ_1/μ_2	2.9	0.31
ν_1	0.25	0.25
ν_2	0.2	0.3

Case 1 corresponds to a hard lining (e.g., high-strength concrete lining in stiff rock masses), whereas Case 2 represents a soft lining (e.g., yielding shotcrete support in weak

or fractured rock). These two scenarios cover typical geotechnical conditions encountered in deep tunnel engineering under blasting loads.

3.2. Asymptotic Degeneration and Validation

To rigorously verify the mathematical consistency and broad numerical robustness of the proposed algebraic matrix framework, a comprehensive degeneration analysis is systematically executed. Prior to the comparative evaluation, a rigorous numerical convergence check of the wave function expansion series was conducted. It was determined that a truncation order of $N = 10$ is fully sufficient to guarantee excellent calculation precision and numerical stability, and is thus universally adopted in this study. To conclusively demonstrate the robustness of the analytical model, we retain a finite explosive source distance of $L/a = 10$ to preserve the pronounced near-field cylindrical wavefront curvature. Concurrently, the tangential frictional slip behavior is completely suppressed ($D^* = 0$), while the dimensionless normal and tangential interface spring stiffness parameters are systematically varied ($K_r^* = K_\theta^* = 10.0, 1.0, 0.1$). Under these specific near-field configurations, the calculated spatial distribution curves of the DSCF for both the surrounding rock and the tunnel lining are meticulously compared with the established analytical solutions presented by Yi et al. [31]. As illustrated in Figure 2, the spatial distribution patterns and peak magnitudes computed by our proposed model perfectly coincide with the existing literature across the entire range of interface compliances. This excellent agreement comprehensively validates the accuracy, stability, and broad robustness of the adopted wave function expansion method and Graf's addition theorem formulations under complex near-field conditions.

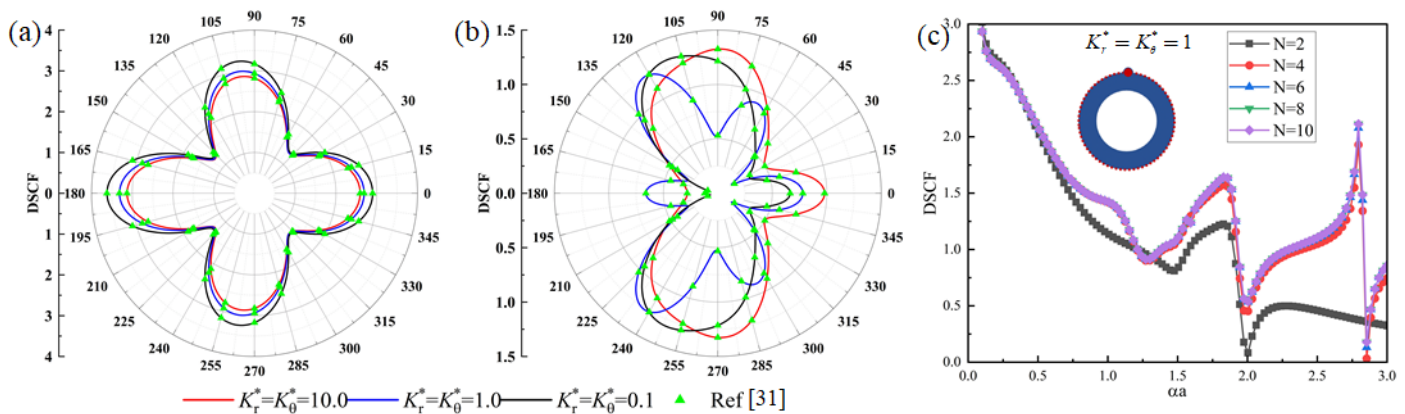


Figure 2. Verification of dynamic response in the rock mass [31] and convergence analysis: (a) $\alpha a = 0.1$; (b) $\alpha a = 2$; (c) convergence of DSCF at $\theta = 90^\circ$ under different truncation orders (N).

4. Influence of the Elastic–Slip Interface and Near-Field Effects on Tunnel Dynamic Response

4.1. Effect of the Elastic–Slip Interface on DSCF

When investigating the impact of the elastic–slip interface on the dynamic stress concentration factor (DSCF), the dimensionless source distance is fixed at $L/a = 10$. To independently analyze the specific mechanisms of K_r^* , K_θ^* , and D^* , a control variable method is adopted: when analyzing K_r^* , we set $K_\theta^* \rightarrow \infty$ and $D^* = 0$; conversely, when analyzing K_θ^* , we set $K_r^* \rightarrow \infty$ and $D^* = 0$. It is important to note that interface compliance and tangential stiffness are physically and mathematically interrelated. Based on the

boundary condition formulations presented earlier, when $K_\theta^* \rightarrow \infty$, the tangential displacement jump satisfies $u_\theta^{(1)} - u_\theta^{(2)} = D\tau_{r\theta}^{(2)}$. Conversely, when $D^* = 0$, this condition transforms into $u_\theta^{(1)} - u_\theta^{(2)} = \tau_{r\theta}^{(2)} / K_\theta^*$. Because these two parameters exert opposing mathematical effects, we temporarily disregard the tangential slip effect (setting $D^* = 0$) when solely discussing the interface stiffness characteristics (K_r^* and K_θ^*) to prevent analytical redundancy.

Figure 3 compares the influence of varying interface spring stiffnesses on the DSCF of the surrounding rock and lining for Cases 1 and 2. It is evident that the dynamic stress in the lining for Case 2 is significantly higher than that in the surrounding rock. This phenomenon occurs because the shear modulus and wave velocity of the lining in Case 2 exceed those of the rock mass, constituting a typical “hard shell–soft matrix” structural system. When an incident wave propagates from the softer rock into the stiffer lining, the pronounced wave impedance mismatch induces intense focusing and amplification effects within the lining, causing its dynamic stress to surge. Furthermore, the lining demonstrates higher sensitivity to changes in spring stiffness than the rock mass; a larger stiffness yields a smoother DSCF distribution curve. A comparison of the subplots in Figure 3 reveals that the response oscillations induced by normal spring stiffness are predominantly concentrated in the low-frequency region, whereas those induced by tangential spring stiffness largely manifest in the high-frequency band. Mechanistically, this occurs because normal stiffness predominantly dictates the transmission of radial compressional waves (characterized by low frequencies), while tangential stiffness governs the transfer of tangential shear waves (characterized by high frequencies). Focusing on low-frequency conditions, Figure 3b,d extract the evolution of the DSCF peak at $\theta = 90^\circ$ under various spring stiffnesses K^* and slip coefficients D^* (with incident wavenumber $\alpha a = 0.1$ and $L/a = 10$). The results indicate that the DSCF evolution exhibits an inverse relationship with K^* and D^* for both cases: a smaller K^* coupled with a larger D^* enhances the energy dissipation and isolation capacity of the interface, thereby reducing the DSCF in the lining.

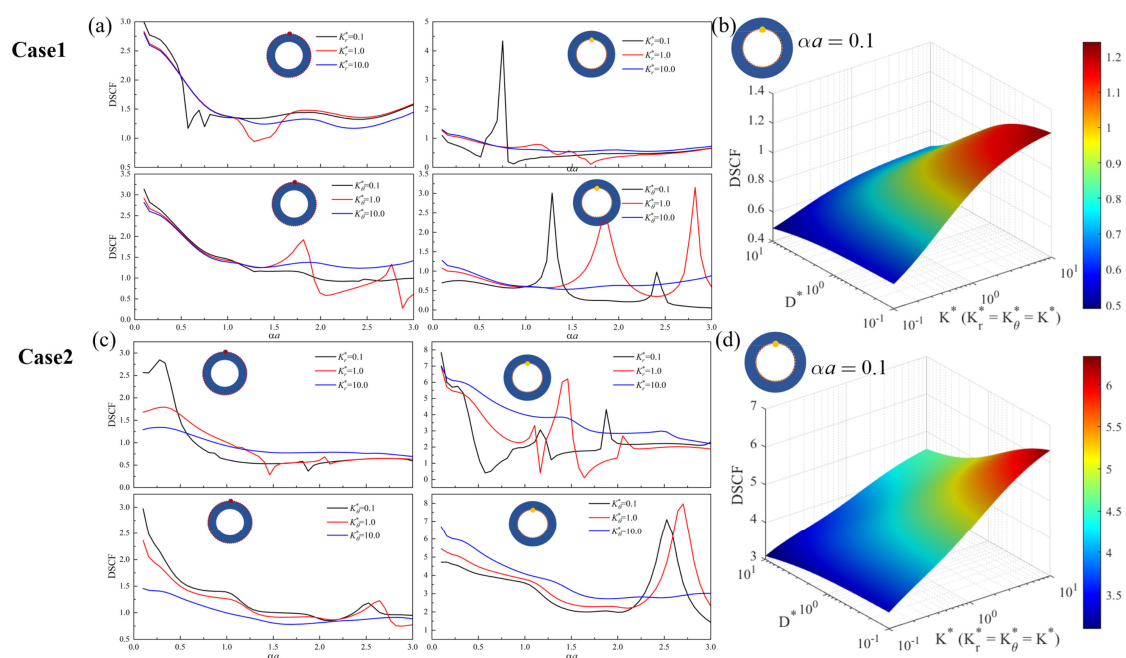


Figure 3. Effect of varying spring stiffnesses on the DSCF of the surrounding rock and lining: (a,b) Case 1; (c,d) Case 2.

To further elucidate the specific impact of the slip coefficient D^* on the DSCF, parameters are set to $K_r^* = K_\theta^* = 1$ and $L/a = 10$. Figure 4 compares the circumferential distribution of the lining DSCF under different slip coefficients for low-frequency ($\alpha a = 0.1$) and high-frequency ($\alpha a = 0.8$) incident waves. The results demonstrate that under low-frequency excitation, the DSCF curves universally exhibit a typical “petal-like” distribution pattern, and the DSCF progressively decreases as the slip coefficient increases, indicating that the interface serves as an effective shock absorber. However, under high-frequency excitation, the response on the blast-facing side of the lining reverses: a larger slip coefficient induces a higher localized DSCF. This suggests that during dynamic disturbances rich in high-frequency components, such as blasting operations, tangential slip effects can provoke more severe stress concentration phenomena. Consequently, the potential negative impact of the slip coefficient must not be overlooked in engineering design.

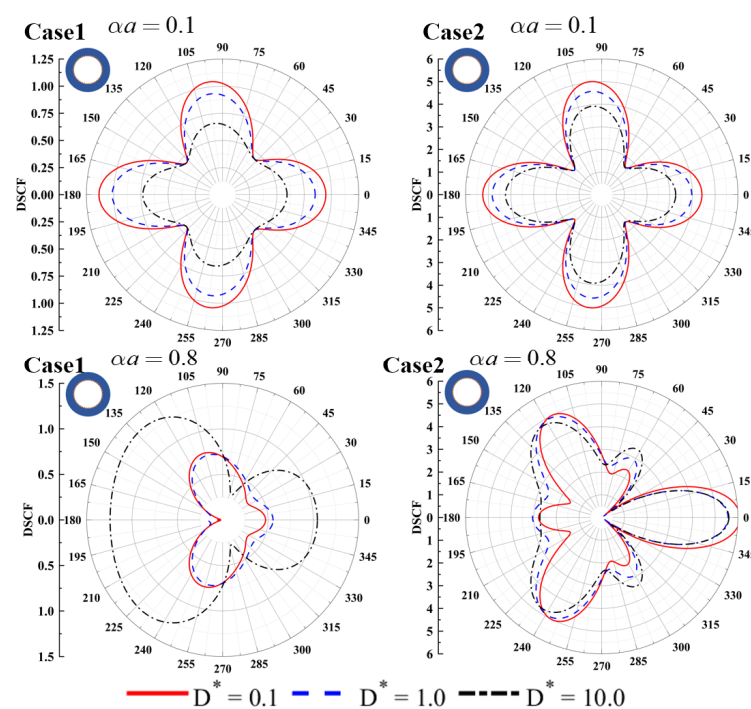


Figure 4. Circumferential distribution of the lining DSCF under varying slip coefficients.

4.2. Effect of Source Distance (Near-Field Effect) on DSCF

Given that the preceding analyses demonstrate that dynamic stress in the lining generally exceeds that in the surrounding rock, this section focuses exclusively on the structural dynamic response of the lining. Considering the coupled effects of the elastic-slip interface, the dimensionless parameters are set to $K^* = 1$ and $D^* = 1$. Figure 5 presents the circumferential distribution of the lining DSCF under varying source distances for both cases, subjected to low- and high-frequency excitations. The graphs clearly illustrate that on the blast-facing side of the lining, a closer source distance results in a larger peak dynamic stress. This occurs because the pronounced curvature of the near-field wavefront induces severe stress concentration and amplification on the blast-facing surface. This compellingly demonstrates that persisting with traditional plane-wave assumptions in near-field blasting analyses will severely underestimate actual structural hazards. As the dimensionless source distance L/a extends to 50 and 100, the DSCF distribution profiles and magnitudes essentially converge, allowing us to conclude that when $L/a \geq 50$, the near-field wave characteristics largely dissipate, justifying a simplification to a plane-wave model. Additionally, comparing the low- and high-frequency responses reveals that

the overall DSCF peak under high-frequency waves is slightly lower than that under low frequencies, but its spatial distribution pattern is markedly more complex. Fundamentally, the short-wavelength characteristics of high-frequency waves exacerbate multiple scattering and interference phenomena between the dual-layer media, highly corroborating the earlier mechanistic analysis. Finally, a comparison between the two lining cases reaffirms that a stiff lining significantly amplifies the dynamic response.

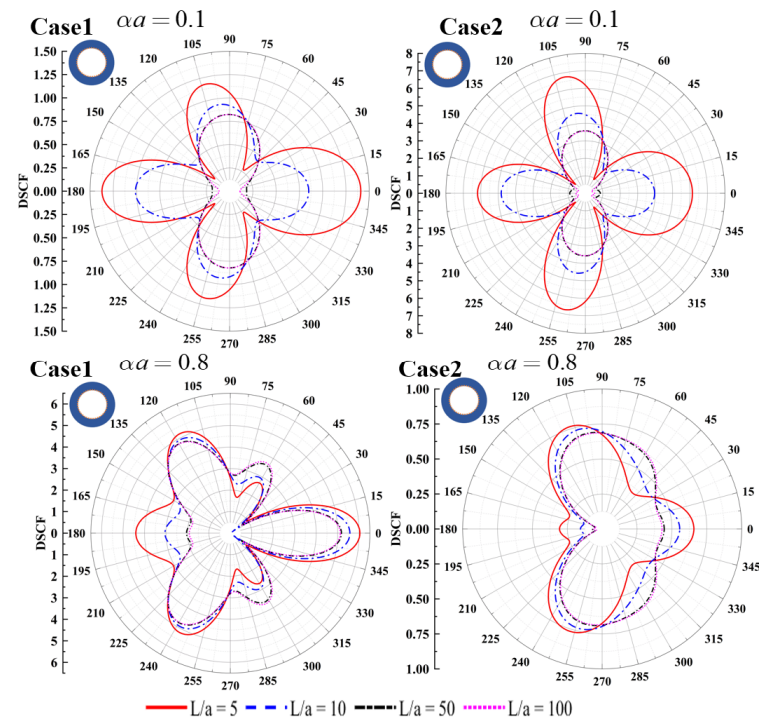


Figure 5. Effect of varying source distances on the circumferential distribution of the lining DSCF.

4.3. Sensitivity Analysis of Interface Mechanical Properties

This section specifically examines the wave characteristics under high-frequency excitation conditions, setting the dimensionless source distance to $L/a = 10$. To vividly unveil the spatial distribution and evolution mechanisms of the wavefield within the media, Figure 6 displays the dynamic stress field distribution contour maps for the two cases under various elastic-slip configurations. It is crucial to note that, to authentically capture the alternating tension and compression interference characteristics exerted by the stress waves, the dynamic stress concentration factor (DSCF) in these contour maps retains its real sign (absolute values are not taken). Observing the contour maps indicates that the propagation patterns of stress waves in the surrounding rock remain fundamentally consistent across different interface conditions, with wavefield variations primarily concentrated within the lining region. Comparing Figure 6a,c reveals that when the interface is perfectly bonded, introducing a slip mechanism (a larger slip coefficient) elevates the peak DSCF, which is consistent with the aforementioned high-frequency trends. Overall, under high-frequency impact, the presence of an imperfect interface generally intensifies the degree of dynamic stress concentration. An in-depth comparison of the two lining cases uncovers an important anomaly: in the stiff lining (Case 2), the imperfect elastic-slip interface effectively mitigates internal stress concentration; however, in the soft lining (Case 1), the exact opposite occurs. Mechanistically, within the present analytical model, the imperfect interface appears to function as a “flexible barrier layer.” For a stiff lining, the results suggest that this barrier obstructs and reflects incoming high-frequency energy, potentially exerting a protective function akin to “isolation stress release” (stronger slip cor-

relates with weaker stress concentration). Conversely, for a soft lining under similar loading conditions, the presence of this barrier may further undermine boundary confinement, leading the lining to undergo intensified localized deformation under dynamic loading, thereby potentially triggering a detrimental “instability stress amplification” effect.

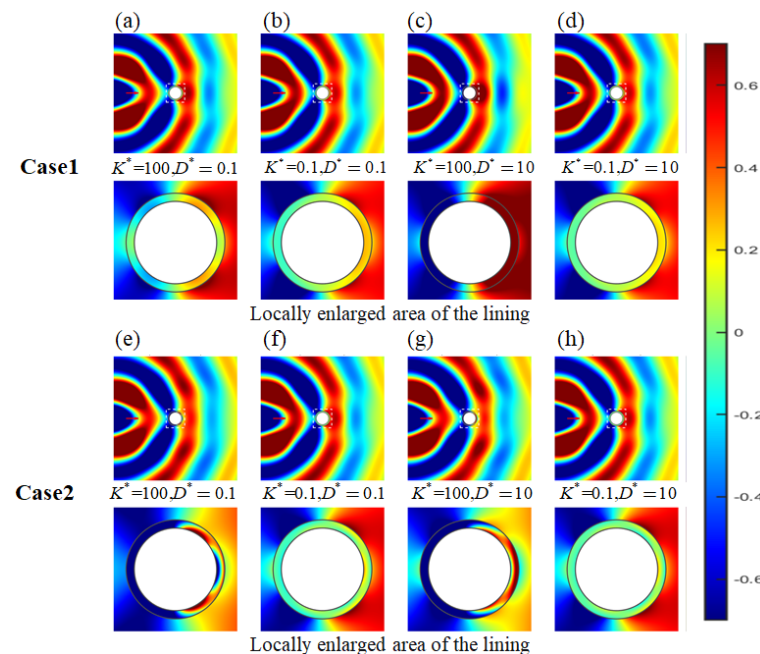


Figure 6. Dynamic stress-field distribution contour-maps under different elastic-slip coefficients: (a–d) Case 1 with $(K^*, D^*) = (100, 0.1), (0.1, 0.1), (100, 10), (0.1, 10)$; (e–h) Case 2 with the same sequence.

4.4. Sensitivity Analysis of Source Distance (Near-Field Effects)

To provide a panoramic visualization of stress transfer and alternating tension–compression interference processes within the rock–lining system, Figure 7 plots the spatial distribution contour maps of the real part of the circumferential dynamic stress concentration factor under high-frequency incidence combined with an imperfect elastic–slip interface.

From the perspective of global wavefield characteristics, in the near-field region, the pronounced wavefront curvature fundamentally alters the spatial topology of the scattered wavefield. When the source distance is relatively short, intense geometric divergence of cylindrical waves combined with structural shadow effects causes the DSCF on the blast-facing side to drastically exceed that on the shadow side. However, as the source distance increases and approaches the far-field limit, the cylindrical wavefront gradually flattens and effectively degenerates into a plane wavefield. At this point, the macroscopic stress field transitions into a typical parallel layered propagation mode, featuring highly symmetric crests and troughs.

Focusing on localized stress responses within the lining clearly highlights the fundamental differences in dynamic behavior between the soft lining (Case 1) and the stiff lining (Case 2). Under identical elastic–slip interface configurations, the overall DSCF peak values within the stiff lining are consistently higher than those in the soft lining, demonstrating that high-stiffness structures are inherently more prone to severe stress accumulation under dynamic loading. More crucially, the high-frequency response contour maps intuitively capture the localized displacement continuity jumps triggered by the imperfect elastic–slip interface. For Case 2 (stiff lining), the elastic–slip interface successfully intercepts and reflects a portion of the incoming high-frequency energy, permitting a degree

of “stress release” within the rigid lining. In stark contrast, for Case 1 (soft lining), the introduction of the imperfect interface further strips away the effective external confinement provided by the surrounding rock, which, under high-frequency impact, paradoxically accelerates the alternating concentration of dynamic tensile and compressive stresses within the soft matrix.

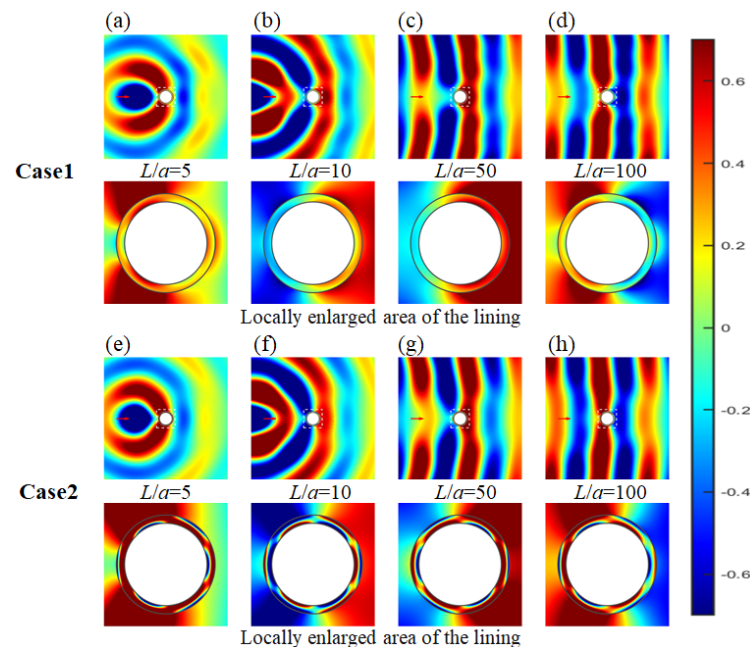


Figure 7. Dynamic stress field distribution contour maps under different source distances: (a–d) Case 1 with $(K^*, D^*) = (100, 0.1), (0.1, 0.1), (100, 10), (0.1, 10)$; (e–h) Case 2 with the same sequence.

5. Conclusions

In this study, the dynamic stress concentration around a circular lined tunnel subjected to cylindrical P-waves is analytically investigated using a generalized elastic–slip interface model. The primary conclusions are drawn as follows:

(1) The curvature of near-field cylindrical waves causes severe energy focusing, significantly amplifying the dynamic stress concentration on the illuminated side of the lining. Traditional plane-wave assumptions underestimate this dynamic risk and should only be applied when the dimensionless source distance exceeds $L/a \geq 50$.

(2) For the parameter range investigated in this study, the imperfect elastic–slip interface exhibits characteristics of a “flexible barrier” whose physical effect is highly sensitive to the relative rock–lining impedance matching. In stiff linings, it suggests a potential to serve as a seismic isolator that reduces stress concentration by providing dynamic stress relief, whereas in soft linings, interfacial slip tends to weaken boundary confinement, triggering unstable local deformations and exacerbating stress concentration.

(3) The interface flexibilities exhibit distinct frequency-dependent behaviors, where normal stiffness predominantly governs low-frequency radial responses, while tangential stiffness and interfacial slip dictate high-frequency behaviors. Neglecting the interfacial slip coefficient under high-frequency dynamic loading (e.g., blasting) leads to inaccurate predictions, severely underestimating the structural vulnerability of soft lining configurations.

Finally, by assuming linear elastic, single-phase media, the present model neglects extreme near-source non-linearities (e.g., rock crushing and plastification) and multiphase groundwater effects. Consequently, our findings apply strictly to the “elastic wave prop-

agation zone” beyond the blast-induced fracture zone. Future work will integrate localized damage and poroelasticity models to capture these complex near-field dynamics.

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Conflicts of Interest: Author Yao Rong and Junping Yu were employed by the Jiangxi Communications Investment Maintenance Technology Group Co., Ltd. The remaining authors declare that the research was conducted in the absence of any commercial or financial relationships that could be construed as a potential conflict of interest.

Appendix A

$$\begin{aligned}
 \varepsilon_{11}^{(l)}(\alpha r) &= \left(n^2 + n - 1/2\beta^2 r^2\right) Q_n^{(l)}(\alpha r) - \alpha r Q_{n-1}^{(l)}(\alpha r) \\
 \varepsilon_{12}^{(l)}(\beta r) &= n \left[-(n+1) Q_n^{(l)}(\beta r) + \beta r Q_{n-1}^{(l)}(\beta r)\right] \\
 \varepsilon_{21}^{(l)}(\alpha r) &= -\left(n^2 + n + 1/2\beta^2 r^2 - \alpha^2 r^2\right) Q_n^{(l)}(\alpha r) + \alpha r Q_{n-1}^{(l)}(\alpha r) \\
 \varepsilon_{22}^{(l)}(\beta r) &= n \left[(n+1) Q_n^{(l)}(\beta r) - \beta r Q_{n-1}^{(l)}(\beta r)\right] \\
 \varepsilon_{41}^{(l)}(\alpha r) &= -n \left[-(n+1) Q_n^{(l)}(\alpha r) + \alpha r Q_{n-1}^{(l)}(\alpha r)\right] \\
 \varepsilon_{42}^{(l)}(\beta r) &= -\left(-n^2 + n - 1/2\beta^2 r^2\right) Q_n^{(l)}(\beta r) + \beta r Q_{n-1}^{(l)}(\beta r) \\
 \varepsilon_{71}^{(l)}(\alpha r) &= \alpha r Q_{n-1}^{(l)}(\alpha r) - n Q_n^{(l)}(\alpha r) \\
 \varepsilon_{72}^{(l)}(\beta r) &= n Q_n^{(l)}(\beta r) \\
 \varepsilon_{81}^{(l)}(\alpha r) &= -n Q_n^{(l)}(\alpha r) \\
 \varepsilon_{82}^{(l)}(\beta r) &= -\left[\beta r Q_{n-1}^{(l)}(\beta r) - n Q_n^{(l)}(\beta r)\right] \\
 Q_n^{(l)}(\cdot) &= \begin{cases} J_n(\cdot), l=1 \\ H_n^{(1)}(\cdot), l=3 \\ H_n^{(2)}(\cdot), l=4 \end{cases}
 \end{aligned}$$

The aforementioned equations detail the calculation factors for the matrix components in Equation (13).

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