HHH-2004: FIRST CARE–HHH–APD WORKSHOP ON BEAM DYNAMICS IN FUTURE HADRON COLLIDERS AND RAPIDLY CYCLING HIGH-INTENSITY SYNCHROTRONS

CERN, Geneva
8–10 November 2004

PROCEEDINGS
Editors: F. Ruggiero, W. Scandale, F. Zimmermann
This report contains the Proceedings of the First CARE–HHH–APD Workshop on Beam Dynamics in Future Hadron Colliders and Rapidly Cycling Synchrotrons, ‘HHH-2004’, held at CERN from 8 to 11 November 2004. The workshop addressed the various ways of boosting the luminosity of the Large Hadron Collider (LHC) by up to a factor $10^{35}$ cm$^{-2}$s$^{-1}$ after about seven years of operation, i.e., around 2014.

The HHH-2004 workshop brought together numerous international experts on hadron colliders, intense proton beams, and accelerator-physics simulations, including numerous particle physicists and some specialists in lepton accelerators, who discussed the LHC and GSI complex upgrade; beam dynamics and intensity challenges; LHC interaction-region upgrade and beam choices; simulation code benchmarking and code repository — for optics, beam–beam effects, conventional impedance, electron cloud and collective instabilities; fast cycling injectors. The simulation codes were discussed in the form of five panel discussions. A poster session with a wide range of contributions provided a welcome opportunity to find further details.

The workshop fostered a fruitful exchange of ideas between various different communities, such as high-intensity proton accelerators, high-energy proton ring colliders, lepton colliders, and the LHC particle-physics experiments. Several joint inter-laboratory simulations and code benchmarking studies were carried out, either in preparation for or inspired by the workshop. As just one example, the mechanism responsible for space-charge-induced halo growth in the GSI FAIR storage rings was invoked by a CERN–GSI collaboration to explain emittance growth due to electron cloud observed in LHC simulations. Topical highlights of the workshop include the studies of IR layouts for a higher-luminosity LHC, the various challenges encountered when raising the LHC beam intensity beyond the design value, and the possible uses of crab cavities, advanced beam–beam compensation techniques, or induction acceleration. The superbunch option for an LHC upgrade could be eliminated at this workshop, since it was found to pose unsolvable problems for the detectors, the beam dump, and the collimation system.

The simulation codes are evolving into two different directions, aiming either for more simplicity or for more complete, self-consistent modelling. Lists of priorities and future tasks for the different types of simulation codes were established, and several workshop participants sketched the prospect of an ultimate universal code. A condensed overview of the primary achievements is provided by the summary reports at the end of these proceedings.
The first CARE–HHH–APD Workshop ‘HHH-2004’, focusing on Beam Dynamics in Future Hadron Colliders and Rapidly Cycling High-Intensity Synchrotrons, was held at CERN from 8 to 11 November 2004 (see http://care-hhh.web.cern.ch/CARE-HHH/HHH–2004). It was attended by more than 100 accelerator and particle physicists, mostly from Europe, but also from the US (LARP Collaboration) and from Japan (KEK). The workshop addressed the various ways of boosting the luminosity of the Large Hadron Collider (LHC) by up to a factor $10^{30}$ cm$^{-2}$s$^{-1}$ after about seven years of operation, i.e., around 2014. Such an LHC upgrade would include new interaction regions, with either stronger or larger-aperture low-beta quadrupoles, moderate modifications of several subsystems like beam dump or collimation, in order to cope with a higher beam current, and, possibly, either long-range beam–beam compensation or crab cavities. The layout of the new interaction region, for which a number of options were presented, is closely linked to both magnet design and beam-dynamics considerations. In a more challenging approach, the upgrade of the LHC injector chain and, in a later stage, the energy upgrade of the LHC itself were envisioned. An LHC injector-chain upgrade could employ concepts similar to those being developed for the GSI FAIR project, the challenges of which were also discussed in great detail at this workshop. A key outcome of the HHH-2004 Workshop is the elimination of the ‘superbunch’ scheme for the LHC luminosity upgrade, since this option was determined to pose unsolvable problems to the detectors, the beam dump, and the collimator system, due to the associated dramatic increase in local charge density. Another important result of the workshop was that it brought together and stimulated the discussion and information exchange between various different communities. In particular, links between LHC accelerator physicists and LHC experimenters were strongly reinforced. They are essential for an overall optimization of the future LHC parameters.

At HHH-2004, a broad range of accelerator-physics issues and beam-dynamics topics relevant to high-intensity high-energy hadron beams were addressed. Recent results of studies on space charge (joint GSI–CERN experiments and simulations), beam–beam effects, electron cloud, and conventional instabilities were presented. The state of the art in simulation tools was revisited and the future direction determined in a combination of overview talks and panel discussions on single-particle codes, beam–beam codes, impedance codes, electron-cloud codes, and instabilities. Most sessions established lists of existing simulation programs and their respective features. It was emphasized that simulation codes should support all stages of an accelerator project with shifting requirements. Often, communication with other specialized codes is required to address specific questions. In response to these observations, an outcome of the discussion was the strong recommendation to develop toolkits and to maintain a modular code structure as well as a standard input format. The latter will also ease code comparison. An example of a common input format is SXF for optics codes, which could be improved in the direction of XML or ADXF. Another possible choice is to introduce a MAD-like input format for other programs, such as for 3D electron-cloud simulation codes. A first example of a modular code structure and also of a new style of code management is the MADX program. The modern codes should ideally provide the possibility of real-time simulations, which will involve judicious compromises between accuracy and speed. Two tendencies, towards more complete, detailed, all-inclusive descriptions on the one hand and towards fast, simplified, few-parameter models on the other were noticed in most of the applications. These two tendencies could be reconciled and supported by a modular code structure. A list of priorities and future tasks for the different types of simulation codes was established, and, in view of rapidly growing computing power, several workshop participants sketched the prospect of an ultimate universal code.

The proceedings are structured according to the seven workshop sessions:

- Session 1: Opening Talks, LHC and GSI Complex Upgrades  
  (chair G. Altarelli, secretary E. Tsesmelis)

- Session 2: Beam Dynamics and Intensity Challenges  
  (chair P. Lebrun, secretary F. Zimmermann)

- Session 3: LHC IR Upgrade and Beam Choices  
  (chair S. Peggs, secretary O. Brüning)

- Session 4: Simulation Code Benchmarking and Repository — A: Optics, Beam–Beam and Conventional Impedance  
  (chairs O. Brüning, W. Herr, T. Weiland)

- Session 5: Fast Cycling Injectors  
  (chair W. Scandale, secretary L. Bottura)

- Session 6: Simulation Code Benchmarking and Repository — B: Electron-Cloud and Collective Instabilities  
  (chairs F. Zimmermann, V.G. Vaccaro)

- Session 7: Workshop Summaries and Panel Discussion on CARE-HH Network  
  (chairs W. Scandale, F. Zimmermann).

These proceedings have been published in paper and electronic form. The paper copy is in black and white; the electronic version contains colour pictures. Electronic copies can be retrieved through: http://care-hhh.web.cern.ch/CARE-HHH/HHH-2004/Proceedings/proceedings_hhh2004.htm

The compilation of these proceedings would not have been possible without the help of the chairmen, scientific secretaries, speakers of the sessions, members of the panel discussions, and presenters of posters. In particular, we would like to thank all the participants for their stimulating contributions.

The HHH-2004 workshop was sponsored and supported by the European Community Research Infrastructure Activity under the FP6 ‘Structuring the European Research Area’ programme (CARE, contract number RII3-CT-2003-506395).

Geneva, 22 May 2005

F. Ruggiero, W. Scandale, F. Zimmermann
CONTENTS

Preface ........................................................................................................................................... v

SESSION 1: OPENING TALKS, LHC AND GSI COMPLEX UPGRADES
 (chair G. Altarelli, secretary E. Tsesmelis)

Possible Scenarios for an LHC Upgrade
F. Ruggiero, F. Zimmermann ................................................................. 1
Introduction to CARE and CERN Strategy 1 R. Aymar
The GSI Future Project 1 W. Henning
Overview of Technological Challenges 1 W. Scandale
US-LARP Contribution to CARE-HHH-APD 1 S. Peggs
LHC Luminosity Upgrade: Physics Motivation and Detector Performance 1 D. Denegri
Relevance of Possible Upgrade of the CERN Accelerator Complex for Fixed Target Physics 1
J. Engelen

SESSION 2: BEAM DYNAMICS AND INTENSITY CHALLENGES
 (chair P. Lebrun, secretary F. Zimmermann)

Machine Protection
K.H. Mess ......................................................................................................................... 15

Vacuum Issues for an LHC Upgrade
O. Gröbner ....................................................................................................................... 19

Space Charge and Optics Studies for High-Intensity Machine
G. Franchetti ................................................................................................................... 25

RF and Feedback for Bunch Shortening
J. Tückmantel ................................................................................................................ 33

Electron Cloud Effects — Observations, Mitigation Measures and Challenges in RHIC and SNS
J. Wei ........................................................................................................................... 43

Optics Design for a Nonlinear Collimation System in the LHC
J. Resta Lopez, A. Faus-Golfe, F. Zimmermann ....................................................... 57

What is an Acceptable Vacuum Pressure in the LHC Arcs?
B. Jeanneret, F. Zimmermann ..................................................................................... 61

Minimum Bunch Length at the LHC
E. Vogel ......................................................................................................................... 67

Morphological and Structural Studies of Crystals for Channeling of Relativistic Particles
S. Baricordi, V. Guidi, C. Malagù, G. Martinelli, E. Milan, M. Stefancich, A. Carnera, A. Sambo, C. Scian,

Limitations due to Electron-Cloud Heat Load for the LHC and Its Upgrade
D. Schulte, F. Zimmermann ....................................................................................... 75
Collimation 1 R. Assmann
Beam Dump 1 B. Goddard

1Paper not submitted to the proceedings. For your convenience, the slides presented are available in electronic form at the URL http://care-hhh.web.cern.ch/CARE―HHH/HHH-2004.
SESSION 3: LHC IR UPGRADE AND BEAM CHOICES

(chair S. Peggs, secretary O. Brünning)

Overview of Possible LHC IR Upgrade Layouts
T. Sen, J. Strait, N. Mokhov ................................................................. 83

First Result of Induction Acceleration in the KEK Proton Synchrotron
K. Takayama .................................................................................. 91

Study of Crab Cavity Option for LHC
K. Ohmi .......................................................................................... 97

Beam–Beam Compensation Schemes
F. Zimmermann ........................................................................... 101

Generation and Benefits of Long Super-bunches
H. Damerau, R. Garoby 121

Machine-Detector Interface and Event Pile-Up: Super-Bunches Vs. Normal Bunches
S. Tapprogge .................................................................................. 131

Beam Dynamics Requirements
O. Brünning

SESSION 4: SIMULATION CODE BENCHMARKING AND REPOSITORY —

A: OPTICS, BEAM–BEAM AND CONVENTIONAL INSTABILITIES

(chairs O. Brünning, W. Herr, T. Weiland)

A Modern Answer in Matter of Precision Tracking: Stepwise Ray-tracing
F. Méot, F. Lemuet ........................................................................ 137

Impedance Calculation and Verification in Storage Rings
K. Bane, K. Oide, M. Zobov ............................................................... 143

Beam–Beam Simulations of Hadron Colliders
T. Sen ............................................................................................ 159

Panel Discussion on Beam–Beam Simulation Codes
M. Furman ...................................................................................... 173

RF Coupling Impedance Measurements Vs. Simulations
A. Mostacci, L. Palumbo, B. Spataro, F. Caspers ................................. 175

Feasibility of Directly Computing Wakes of Lossy Collimators
W. Bruns ......................................................................................... 183

Cold-to-Warm Transition Impedances in the SPS Machine
B. Spataro, D. Alesini, M. Migliorati, A. Mostacci, L. Palumbo, F. Ruggiero ............................................................... 185

Coherent Beam–Beam Modes in the LHC for Multiple Bunches, Different Collision Schemes, and Machine Symmetries
W. Herr, T. Pieloni .......................................................................... 195

Overview of Single Particle Codes
W. Decking

1Paper not submitted to the proceedings. For your convenience, the slides presented are available in electronic form at the URL http://care-hhh.web.cern.ch/CARE-HHH/HHH-2004.
SESSION 5: FAST CYCLING INJECTORS  
(chair W. Scandale, secretary L. Bottura)

High Brightness Protons Beams for LHC  
R. Garoby, M. Benedikt .................................................. 203

SPS Impedance and Intensity Limitations  
A. Shaposhnikova .......................................................... 209

Multi-Turn Extraction Based on Trapping in Stable Islands  
R. Cappi, S. Gilardoni, M. Giovannozzi, M. Martini, E. Métal, J. Morel, P. Scaramuzza, R. Steerenberg, A.-S. Müller .......................................................... 215

Fixed Field Alternating Gradient Synchrotrons  
F. Méot .................................................. 221

SIS100/300 & High Energy Beam Transport¹  
P. Spiller

Fast Pulsed SC Magnets for SIS and Super-SPS¹  
G. Moritz

SESSION 6: SIMULATION CODE BENCHMARKING AND REPOSITORY —  
B: ELECTRON-CLOUD AND COLLECTIVE INSTABILITIES  
(chairs F. Zimmermann, V.G. Vaccaro)

Overview of Single-Beam Collective Instabilities in Circular Accelerators  
E. Métal .................................................. 227

Intensity Limitations by Combined and/or (Un)Conventional Impedance Sources  
G. Rumolo .................................................. 249

Code Comparisons and Benchmarking with Different SEY Models in Electron-Cloud Build-Up Simulations  
G. Bellodi .................................................. 261

Electron-Cloud Instability Simulations for the LHC  
E. Benedetto, F. Zimmermann .................................................. 267

Coherent Radiation of Electron Cloud  
S. Heifets .................................................. 271

Overview of Electron-Cloud Simulation Codes¹  
M. Furman

¹Paper not submitted to the proceedings. For your convenience, the slides presented are available in electronic form at the URL http://care-hhh.web.cern.ch/CARE-HHH/HHH–2004.
SESSION 7: WORKSHOP SUMMARIES
(chairs W. Scandale, F. Zimmermann)

Summary of Session 1: LHC and GSI Complex Upgrades
E. Tsesmelis ................................................................. 275

Summary of Session 2: Beam Dynamics and Intensity Challenges
P. Lebrun, F. Zimmermann ........................................... 279

Summary of Session 3: LHC IR Upgrade and Beam Choices
S. Peggs, O. Brüning ..................................................... 283

Summary of Session 5: Fast Cycling Injectors
W. Scandale ................................................................. 287

Summary of Panel Discussion 1: Single Particle Codes
O. Brüning ................................................................. 289

Summary of Panel Discussion 2: Beam–Beam Codes and Simulations
W. Herr ................................................................. 291

Summary of Panel Discussion 3: Impedance Codes
T. Weiland ................................................................. 295

Summary of Panel Discussion 4: Electron Cloud Codes
F. Zimmermann .......................................................... 297

Summary of Panel Discussion 5: Coherent Instabilities
V.G. Vaccaro, G. Rumolo ............................................. 301

List of Participants ....................................................... 305
Possible Scenarios for an LHC Upgrade

F. Ruggiero, F. Zimmermann, CERN, Geneva, Switzerland

Abstract
After discussing the rationale for an LHC upgrade and various performance limitations of the nominal LHC, several alternative upgrade scenarios and options are presented.

1 THE NEED FOR AN UPGRADE

There are two reasons for considering an upgrade of the LHC after about 6 years of operation. They are illustrated in Fig. 1, which goes back to J. Strait [1]. The figure illustrates two possible evolutions of peak luminosity as a function of the year, as well as the corresponding integrated luminosity and the run-time needed to halve the statistical error of experimental measurements. Both scenarios assume an LHC start up in 2007, and that 10% of the design luminosity is reached in 2008, and 100% in 2011 [2, 3]. In one case, the luminosity is taken to be constant from then on, in the other it continues to increase linearly until the so-called ultimate luminosity (2.3 times the nominal) would be reached by 2016. The radiation damage limit of the LHC low-β quadrupoles is estimated at an integrated luminosity of 600–700 fb⁻¹ [4]. As the figure shows, this value would be exceeded in 2014 or 2016 depending on the scenario. The additional run-time required to halve the statistical error rises more steeply. It would exceed 7 years by 2011 or 2013, respectively.

Since the life expectancy of the interaction-region (IR) magnets is, therefore, estimated to be less than 10 years due to the high radiation doses and since the time needed to halve the statistical error will exceed 5 years by 2011–2012, it is reasonable to plan a machine luminosity upgrade based on new low-β IR magnets before about 2014.

2 CHRONOLOGY OF UPGRADE STUDIES

In the summer of 2001 two CERN task forces investigated the physics potential [6] and accelerator aspects [5] of an LHC upgrade. Also in 2001 T. Sen, J. Strait, and A. Zlobin published a first paper on the LHC IR upgrade [7]. It was followed by several others from US and CERN teams [8, 9, 10, 2]. In March 2002 an LHC IR upgrade collaboration meeting was held at CERN [11], followed in October 2002 by an ICFA seminar on ‘Future Perspectives in High Energy Physics’. Also the first LHC performance workshop, Chamonix 2003, addressed questions relevant for an upgrade [12]. Since 2004, as part of the 6th framework programme of the European Union the CARE-HHH European Network on High-Energy High-Intensity Hadron Beams [13] is pursuing the LHC upgrade. This network has hosted the HHH-2004 workshop.

Figure 1: Time to halve the statistical error (green curves), integrated luminosity (blue curves), and peak luminosity (red curves) for two different scenarios compatible with the baseline LHC: (1) the luminosity is raised to the nominal value by 2011 and then stays constant, (2) it continues to increase linearly, reaching the ultimate value by 2016. The assumed radiation damage limit of the IR magnets is 700 fb⁻¹ [J. Strait [11]].

3 PERFORMANCE LIMITATIONS

Past studies [5]. have identified a number of LHC performance limitations.

The existing beam dumping system limits the total current and may require an upgrade. The present system is compatible with the ultimate bunch intensity of 1.7 × 10¹¹ protons per bunch. Increases to 2.0 × 10¹¹ could be tolerated with reduced safety margin or after a moderate upgrade.

The detector architecture also limits the luminosity. A detector upgrade can take place in parallel with the accelerator upgrade. It could allow moving the low-β quadrupoles closer to the IP. In their present configuration, the ATLAS and CMS detectors can accept a maximum luminosity of 3–5 × 10³⁴ cm⁻²s⁻¹.

The collimation system and machine protection restrict the total current and the minimum achievable β⁺ value. The machine protection is challenging already for the nominal LHC: The transverse beam energy density is 1000 times that of the Tevatron. The use of simple graphite collimators may limit the maximum transverse energy density to one half of the nominal value, in order to prevent collimator damage and beam instabilities. Closing collimators to 6σ
yields an impedance at the edge of instability. A local fast loss of \(2.2 \times 10^{-6}\) of the beam intensity quenches nearby arc magnets.

The electron-cloud effect may constrain the minimum bunch spacing. The electron cloud in the LHC arcs gives rise to additional heat load on the beam screen. The value of this heat load depends on the surface parameters. At 75-ns bunch spacing, no problem is anticipated. Initial bunch populations at the nominal 25-ns spacing will be limited to about half of the nominal intensity.

As in previous colliders, the beam-beam effect will limit the ratio of bunch intensity and normalized emittance \(N_b/(\gamma \epsilon)\) and the crossing angle. Beam-beam compensation schemes may help to relax the associated constraints.

In the following we will review two of the limiting effects in more detail, namely the electron cloud and the beam-beam interaction.

### 3.1 Electron Cloud

Figure 2 shows a schematic of the electron build-up process inside the LHC beam pipe. Photo-electrons are created by synchrotron radiation at the beam-pipe wall. These photo-electrons are then accelerated in the electric field of the photon-emitting bunch which passes simultaneously. They gain a maximum energy of 200 eV, close to the energy where the secondary emission yield is maximum, and hit the opposite side of the vacuum chamber after about 5 ns. Upon impact on the wall, the accelerated primary electrons generate low-energy secondary electrons, which may stay inside the beam pipe until the following bunch arrives, 25 ns behind the previous one. The secondary electrons are then accelerated by the field of this bunch, producing new secondary electrons in turn. The repetition of this process leads to an avalanche-like generation of electrons. The build-up of electrons only saturates, when the electron space-charge field prevents further secondary electrons from penetrating into the inside of the vacuum chamber. The survival of electrons between bunches and, hence, the rate and magnitude of the electron build-up are enhanced by a high probability of elastic reflection of low-energy secondary electrons [14], i.e., ones which hit the wall before being accelerated by a passing bunch.

An exponential electron cloud build up, as described above, with saturation after about 30–40 bunches has been observed in the SPS with LHC beam [15]. This demonstrates that, even in the absence of synchrotron radiation and a large number of photoelectrons as primary source, the amplification via acceleration by the beam and subsequent secondary emission is strong enough to lead to a rapid increase in the number of electrons, presumably starting from the small rate of electrons generated by gas ionization. In consequence, an electron-cloud effect is anticipated in the LHC also at injection.

It is expected that after some bombardment with electrons, the maximum secondary emission yield will decrease to values close to 1, whence the electron-cloud effect will disappear. In other words, it is thought that the electron cloud will cure itself (so-called 'scrubbing effect'), which will allow reaching the nominal LHC parameters. Experiments in the SPS have verified the scrubbing phenomenon and have established that the desired low values of secondary emission can be reached.

Nevertheless, it must be emphasized that the electron-cloud build up steeply increases for lower bunch spacing. Only experience with LHC operation will clarify whether bunch spacings lower than the nominal 25 ns are prohibited.

![Figure 2: Schematic of electron-cloud build up in the LHC vacuum chamber. Primary electrons are generated on the chamber wall illuminated by synchrotron radiation via photomission. The number of electrons is then amplified exponentially by beam-induced multipacting.](image-url)
experimental data. In Fig. 4 the bunch population at which an electron-cloud effect is observed at various accelerators is displayed as a function of the bunch spacing [18]. Note that both axes carry a logarithmic scale. Data for positron as well as hadron storage rings are included. For a large class of storage rings the threshold bunch population seems to scale linearly with the bunch spacing:

\[ N_b^{th} \propto \Delta t_{sep}. \] (1)

At a spacing of 20–25 ns, the points for the SPS, PS, Tevatron with uncoalesced beam, and the APS fall on top of each other. For the SPS, data of electron-cloud thresholds are available for three different spacings (5 ns, 25 ns and 50 ns), and follow the empirical scaling of (1). Also RHIC data reveal a large sensitivity to the bunch spacing, but the observed thresholds are much lower than for the other machines. A tentative explanation is that the surfaces in RHIC may be less well conditioned, and, therefore, exhibit a larger secondary emission yield. The thresholds found at DAFNE and KEKB are looser than for the proton machines and the APS, possibly due to extensive surface cleaning by synchrotron radiation. Shown in red are design operating points of several planned accelerators, including the nominal and ultimate LHC. From here the upgrade path could go in two different directions, either parallel to the SPS ‘threshold line’ towards larger spacing and higher intensity with \( N_b \propto \Delta t_{sep} \), staying at a constant distance from the electron-cloud threshold, or towards shorter bunch spacings \( \Delta t_{sep} \) for constant bunch intensity \( N_b \), which would require an additional improvement in the surface conditions.

Figure 5 visualizes the consequences of a shorter bunch spacing for the heat load in LHC, presenting simulation results for two different values of \( \delta_{max} \). The figure shows that the change of \( \delta_{max} \) result in a roughly constant vertical shift on a vertical logarithmic scale. Elastically reflected low-energy electrons were taken into account (here assuming 50% reflection probability in the limit of 0 incident electron energy.

At low and moderate bunch intensities the electron-cloud build up saturates, if the average electron line density equals the average beam charge line density [19]. This leads to an equilibrium electron density which scales linearly with the beam intensity. However, the electron density no longer increases for high bunch intensities [20, 21], or at \( N_b > N_{trans} \), with the critical or transition intensity given by

\[ N_{trans} = \frac{E_s \Delta t_{sep}}{m c e}, \] (2)

where \( E_s \) denotes the typical kinetic energy of the emitted
secondary electrons, $\Delta t_{\text{sep}}$ the bunch spacing in seconds, $m_e$ the electron mass, and $r_e$ the classical electron radius. For the LHC bunch spacing, we have $N_{\text{trans}} \approx 10^{11}$, assuming $E_s \approx 1.9$ eV [22]. The transition occurs, when the initial kinetic energy $E_s$ of the secondary electrons is too low for them to penetrate into the space-charge field of the electron cloud [20].

The saturated electron line density in the high-intensity regime is estimated as

$$N_e^{\text{sat}} \approx \frac{E_s}{r_e m_e c^2} \approx 1.3 \times 10^9 \text{ m}^{-1},$$

where again we have taken $E_s \approx 1.9$ eV as the typical secondary-electron energy (see [22]).

Figure 6 shows simulations which reveal the predicted saturation for high intensity. Shown is the simulated electron line density evolution in time during the passage of four consecutive LHC bunch trains, with bunch populations of $N_b = 2.3 \times 10^{11}$ and $4.6 \times 10^{11}$, respectively. However, the fact that the electron density barely changes for higher intensity does not mean that the heat load remains constant. Since the energy transfer to an electron roughly scales with the square of the bunch charge (if the electron is ‘far’ from the beam), the electron-induced heat load continues to grow for increasing intensities.

Figure 6: Simulated build up of electron-cloud line density (in m$^{-1}$) vs. time (in s) for bunch populations of $N_b = 2.3 \times 10^{11}$ and $4.6 \times 10^{11}$, at the nominal 25-ns bunch spacing. Assumed were a maximum secondary emission yield $\delta_{\text{max}} = 1.3$, the energy for which the yield is maximum $E_{s,\text{max}} = 230$ eV, and a primary photo-electron emission rate $d\lambda_{e,p}/ds = 7.25 \times 10^{-4}$ m$^{-1}$ per proton [18].

Figure 7 illustrates electron motion during the passage of a long superbunch [23] with nearly uniform profile. Photoelectrons generated at the head of the bunch are trapped in the increasing beam potential and released only at the end of the bunch passage. Electrons emitted at the wall during most of the bunch passage move in a quasi-static beam potential, and do not gain any net energy from the beam. They traverse the beam, being first accelerated and then decelerated, and hit the opposite side of the chamber with their original emission energy, which is too low to produce a significant amount of secondary electrons. Only electrons generated near the very tail of the bunch experience a beam potential decreasing in time and, as a result, experience a net energy gain. These electrons can therefore contribute to an amplification process, which is appropriately called ‘trailing-edge multipacting’ [24]. The severity of the trailing-edge multipacting depends on the detailed shape of the bunch profile. In any case, the large majority of protons in a superbunch do not participate in the multipacting process, and, therefore, the heat loads calculated for superbunches tend to be negligible, orders of magnitude below those for the nominal LHC bunched beam [25].

3.2 Beam-Beam Interaction

Figure 8 shows a layout of the LHC double ring with its four main experimental insertions. The beam-beam interaction in LHC consists of two components: quasi-head-on beam-beam collisions at the four primary interaction points (IPs) in the four experiments and a total of 120 long-range collisions experienced when approaching or leaving these IPs. The long-range collisions are illustrated in Fig. 9. Of particular importance are so-called PACMAN bunches at the head or tail of a bunch train [26], which do not suffer the nominal number of long-range collisions, and, in consequence, have orbits and tunes different from the nominal bunches. The long-range collisions perturb the motion at large betatron amplitudes where particles come close to the opposing beam. They thereby may cause a ‘diffusive’ (or dynamic) aperture [27], possibly resulting in high background and poor beam lifetime. The number of long-range collisions has increased from only 9 at the SppS collider, over 70 for the Tevatron, to 120 in the LHC, rendering their effect more and more significant.

The LHC dynamic aperture due to the long-range collisions, normalized to the rms beam size $\sigma$, can be estimated as [28, 29, 5]

$$\frac{d\alpha}{\sigma} \approx \theta \sqrt{\frac{\beta^*}{(\gamma\varepsilon)}} \sim 3 \sqrt{\frac{3.75 \mu m r_{\text{beam}} N_b}{30 \times 10^{11}}},$$  

(4)
where \( n_{\text{par}} \) denotes the total number of parasitic (long-range) collisions around one IP (with a nominal number of 30), \( \theta \) is the full crossing angle, and \( \beta^* \) the beta function at the collision IP. If a certain diffusive aperture is required, (4) determines the minimum acceptable crossing angle. Note that for a smaller \( \beta^* \) or a larger bunch intensity, the crossing angle \( \theta \) must be increased to preserve the same diffusive aperture.

Figures 8 and 9: Layouts of the LHC with its four interaction points [30] and schematic of long-range collisions occurring around an LHC interaction point for nominal parameters, including ‘PACMAN’ bunches at the start and of a train [31].

A variable describing the strength of the head-on beam-beam interaction is the beam-beam tune-shift parameter, which for one IP reads

\[
\xi_{\text{HO}} = \frac{N_0 \sigma_p}{4 \pi \gamma \epsilon}.
\]

If the acceptable beam-beam tune shift \( \xi_{\text{HO}} \) is limited, for example by vicinity to nearby resonances, the beam brilliance \( N_0 / (\gamma \epsilon) \) is also limited via (5).

In a similar way, the tune shift induced by the long-range collisions around one IP can be quantified as [26]

\[
\xi_{LR} = 2n_{\text{par}} \frac{\xi_{\text{HO}}}{d^2},
\]

where \( d \) signifies the normalized separation (in units of the rms beam size \( \sigma \)). The long-range tune shift increases as either the bunch spacing (larger \( n_{\text{par}} \)) or the crossing angle is reduced (smaller \( d \)).

Values of head-on tune-shift parameters for a single IP and the total beam-beam tune spread are compared in Table 1 for various hadron colliders. The LHC design parameters confine the total beam-beam tune spread to 0.01 (including long-range contributions), which, based on the SPS collider experience, is taken to be a conservative value of the tune spread.

Table 1: Head-on beam-beam tune shift parameter per IP, no. of IPs and total head-on beam-beam tune shift for four hadron colliders.

<table>
<thead>
<tr>
<th></th>
<th>( \xi_{\text{HO}} )/IP</th>
<th>no. of IPs</th>
<th>( \Delta Q_{\text{bb}} ) total</th>
</tr>
</thead>
<tbody>
<tr>
<td>SPS</td>
<td>0.005</td>
<td>3</td>
<td>0.015</td>
</tr>
<tr>
<td>Tevatron (pbar)</td>
<td>0.01–0.02</td>
<td>2</td>
<td>0.02–0.04</td>
</tr>
<tr>
<td>RHIC</td>
<td>0.002</td>
<td>4</td>
<td>~0.008</td>
</tr>
<tr>
<td>LHC</td>
<td>0.0034</td>
<td>2 (4)</td>
<td>~0.01</td>
</tr>
</tbody>
</table>

In the case of a superbunch, the transition between head-on and long-range collisions becomes continuous, as is illustrated by Fig. 10. One advantage of a uniform superbunch is that, according to calculations [33, 32, 34], the luminosity with long bunches with flat longitudinal distribution is 1.4 times higher than for conventional Gaussian bunches with the same beam-beam tune shift and identical bunch population.

Figure 10: Schematic of superbunch collisions consisting of head-on and long-range components [32].
4 UPGRADE SCENARIOS AND OPTIONS

The fundamental luminosity equation is

\[
L \approx \frac{n_b N_c^2 f_{\text{rep}}}{4\pi\sigma^*} F, \tag{7}
\]

where

\[
F \approx \frac{1}{\sqrt{1 + \left(\frac{\theta_c}{2\sigma^*}\right)^2}}, \tag{8}
\]

and the rms beam size is related to the IP beta function and emittance via the usual relation

\[
\sigma^* = \sqrt{\beta^* \epsilon}. \tag{9}
\]

From (7) and (8), below the beam-beam limit the luminosity is reduced for long bunches and large crossing angles. Hence, in this regime, one would like to shorten the bunch and reduce \(\theta_c\) to maximize the performance.

The situation changes, if we introduce the limitations from the beam-beam interaction. For simplification, we consider two high-luminosity interaction points with alternating crossing (so that the linear beam-beam tune shift due to the long-range collisions cancels between the two IPs [26]). In this case, the total linear tune shift is reduced by a factor \(F_{bb}\) which is approximately equal to the factor \(F\) in (7) [33]. Namely, we have for the total tune shift

\[
\Delta Q_{bb} = \xi_x,\text{HO} + \xi_y,\text{HO} \approx \frac{N_{\text{IP}} F_{bb}}{2\pi \gamma\epsilon}. \tag{10}
\]

Combining (7) and (10), we may rewrite the luminosity as

\[
L \approx \gamma (\Delta Q_{bb})^2 \frac{\pi (\gamma \epsilon) f_{\text{rep}}}{\gamma_p^2 \beta^*} \sqrt{1 + \left(\frac{\theta_c \sigma^*}{2\sigma^*}\right)^2}. \tag{11}
\]

Equation (11) shows that at the beam-beam limit, the luminosity can be increased by increasing the bunch length or the crossing angle, which is the opposite of the behavior below the beam-beam limit. Closer inspection of (11) also reveals that a higher injection energy, which would allow injecting a more intense beam with a larger normalized emittance (\(\gamma \epsilon\)), would raise the luminosity. Another possibility to achieve higher luminosity is to operate with large crossing angle (either in a regime with large Piwinski angle [33] or, in the extreme case, ‘superbunches’ [23]). Figure 11 illustrates the potential luminosity gain in LHC luminosity vs. bunch length (or crossing angle) for Gaussian and flat (super-)bunches at constant beam-beam tune shift with alternating crossings in IP1 and IP5. A more precise prediction, taking into account the bunch-end effects, was derived by H. Damerau [35].

From (7)–(11) two alternative upgrade scenarios emerge.

4.1 Baseline Scheme

Figure 12 shows a flowchart of the baseline upgrade scheme. Starting with the nominal performance at

\[
0.58 \text{ A beam current and luminosity 1 (in units of } \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}), \text{ the bunch intensity is increased until the beam-beam limit is reached at a luminosity of about 2.3 with only two high-luminosity experiments. Then the form factor } F \text{ is increased. Here, several options exist. Long-range beam-beam compensation [36] allows reducing the crossing angle or squeezing } \beta^* \text{ with a constant } \theta_c. \text{ Once the crossing angle cannot be decreased any further, } F \text{ can be further decreased by shortening the bunches using a higher frequency rf system and a 30% smaller longitudinal emittance.}
\]

A completely different approach would be the deployment of crab cavities. The crab-cavities provide effectively head-on collisions, as concerns luminosity, though the bunch centroids still intersect with a large crossing angle. The crab cavities would additionally permit a substantial increase in the crossing angle, to several mrad, whereby the two beams could be passed through separate magnetic channels, without sharing a quadrupole aperture. The crab cavities would lead to a simplified IR design with large \(\theta_c\).

Independently from the procedure followed for maximizing the form factor \(F\), the heartpiece of all LHC upgrades is the installation of new IR magnets, which can provide at least a factor 2 reduction in the IP beta function, to a value of 0.25 m or smaller. The smaller \(\beta^*\) will double the luminosity to 4.6. In case the constraints from electron cloud, beam dump and impedance are not prohibitive, finally the number of bunches can be increased by a factor of about 2, yielding a total luminosity gain by a factor 9.2 above the nominal with a beam current of 1.72 A.

4.2 Piwinski Scheme

The alternative upgrade path moves in the direction of longer bunches with larger crossing angle. A flowchart for this upgrade scenario is depicted in Fig. 13. Again starting from the nominal LHC parameters, as for the baseline scheme, the beta function is reduced to half the nominal...
value by means of new IR magnets. In this upgrade scheme we enlarge the Piwinski parameter and thus decrease the form factor $F$, by increasing the product of bunch length and crossing angle. To maximize the luminosity gain, the longitudinal profile should be flattened, so that the line density is roughly constant along the full bunch length. At the same time as the bunch is lengthened, the bunch charge must be increased to stay at the beam-beam limit. The total number of bunches can be reduced to limit the total beam current. In the extreme limit only a single superbunch would remain. At an average beam current of 0.86 A a luminosity increase by factor 7.7 above the nominal is possible. For a current of 1.72 A, as considered in the baseline scheme, the luminosity gain would be a factor 15.5, which is more than 50% higher than for the baseline scenario at equal current.

4.3 Additional Considerations

If the total beam current is limited, e.g., by electron cloud, machine protection, or dump, fewer bunches with more charge yield a higher luminosity, but also increase the event pile up in the physics detectors.

The minimum value of the IP beta function depends on the interaction-region magnets, the chromatic correction (more critical for scenarios with larger momentum spread), and the settings of the collimators.

The integrated luminosity scales as $T_{bb}/(T_{bb} + T_{\text{turn-around}})$, i.e., with the ratio of collision time and the total time, which is the sum of collision and turn-around time. The turn-around time can be decreased by increasing the injection energy into the LHC (by a “Super-SPS”), which reduces injection time and snapshot.

The compensation of long-range beam-beam encounters and the Super-SPS would enable operation with larger intensity at larger normalized transverse emittance. Increasing the injection energy by a factor of 2, the LHC luminosity also increases by roughly a factor 2.

As pointed out before, the luminosity at the beam-beam limit is higher for flat (long) bunches, instead of Gaussian ones.

The capability of the experiments, e.g., which bunch structures can be handled, must be taken into account.

4.4 Upgrades to the LHC Injector Complex

A possibility being considered also for the CNGS proton beam (CERN Neutrino beam to Gran Sasso) is to upgrade
Figure 13: Schematic of LHC upgrade for operation in a regime with large Piwinski angle.

Figure 14 displays two ‘baseline’ upgrade schemes. Shown on the left is an option with short bunches colliding at a small crossing angle, facilitated by long-range beam-beam compensation schemes. It also reduces the turn-around time and increases the integrated luminosity. A Super-SPS would mark the first step in the direction of an LHC energy upgrade, since it reduces the energy swing by a factor of 2.

A superconducting linac could replace the PS Booster. Alternatives would be a fast cycling superconducting Super-PS and Super-PSB or a series of FFAGs (fixed-field alternating gradient synchrotrons) as proposed for the BNL site [37].

4.5 New Interaction Regions

The goal of the new interaction regions is to reduce $\beta^*$ by a factor 2–5. Various optics are being considered [7, 8, 9, 10, 2, 41]: a ‘cheap’ upgrade based on NbTi as well as stronger magnets made from NbTi(Ta) or Nb$_3$Sn. Both the low-$\beta$ quadrupoles and the separation dipoles should be upgraded.

Several factors drive the IR design [2]: (1) minimization of $\beta^*$, (2) minimization of long-range collision effects, (3) large radiation power directed towards the interaction regions, (4) accommodation of crab cavities or beam-beam compensators, and (5) compatibility with the upgrade path. Items (1) and (2) are addressed by maximizing the magnet aperture and minimizing the distance to the IP.

Figure 14: Schematic of LHC upgrade for operation in a regime with large Piwinski angle.
cause of the large crossing angle, this scheme must employ either long bunches (Piwinski regime) or operate with crab cavities.

Lastly, in case the LHC luminosity must be doubled earlier than foreseen, a more conventional solution, based on NbTi quadrupoles, does also exist [41]. This solution is shown in Fig. 16. Here, the length and aperture of each quadrupole are individually optimized, and the IP-quadrupole distance is slightly reduced from 23 m to 22 m. A beta function of $\beta^* = 0.25$ m appears possible [41].

4.6 Bunch Structure

Various bunch structures correspond to the different upgrade paths. Figure 17 sketches the possible evolutions.

In the baseline upgrade path, we shorten the bunches and increase their number. Advantages of this path are that crab cavities can be used and the event pile up remains tolerable. Concerns are the electron cloud, long-range beam-beam interaction, and impedance.

In the alternative second upgrade path, the trend is towards longer and fewer bunches (e.g., 75 ns spacing). The merits of this scheme include the absence of an electron cloud and a higher luminosity for equal beam current. Concerns are the event pile up and also the impedance. The extreme version of this upgrade direction would be a superbunch, which has the same advantages, offering even slightly higher luminosity, but which, as a striking disadvantage, implies an enormous number of pile-up events, which cannot be handled by the present detectors or their moderate upgrade.

Transitions between the different bunch structures would be possible in the LHC itself, by employing new rf systems for bunch merging or bunch splitting (see also [35]).
Table 2: Parameters for the nominal and ultimate LHC compared with those for three upgrade scenarios with (1) shorter bunches at 12.5-ns spacing [baseline], (2) longer more intense uniform bunches at 75-ns spacing [large Piwinski parameter], and (3) a single superbunch per ring [very large Piwinski parameter].

<table>
<thead>
<tr>
<th>parameter</th>
<th>symbol</th>
<th>nominal</th>
<th>ultimate</th>
<th>shorter bunches</th>
<th>longer bunches</th>
<th>superbunch</th>
</tr>
</thead>
<tbody>
<tr>
<td>no. bunches</td>
<td>(n_b)</td>
<td>2808</td>
<td>2808</td>
<td>5616</td>
<td>936</td>
<td>1</td>
</tr>
<tr>
<td>protons/bunch</td>
<td>(N_0) [10^{11}]</td>
<td>1.15</td>
<td>1.7</td>
<td>1.7</td>
<td>6.0</td>
<td>5600</td>
</tr>
<tr>
<td>bunch spacing</td>
<td>(\Delta t_{\text{sep}}) [ns]</td>
<td>25</td>
<td>25</td>
<td>12.5</td>
<td>75</td>
<td>89000</td>
</tr>
<tr>
<td>average current</td>
<td>(I) [A]</td>
<td>0.58</td>
<td>0.86</td>
<td>1.72</td>
<td>1.0</td>
<td>1.0</td>
</tr>
<tr>
<td>norm. transv. emittance</td>
<td>(\gamma_r) [\mu m]</td>
<td>3.75</td>
<td>3.75</td>
<td>3.75</td>
<td>3.75</td>
<td>3.75</td>
</tr>
<tr>
<td>longit. profile</td>
<td></td>
<td>Gaussian</td>
<td>Gaussian</td>
<td>Gaussian</td>
<td>uniform</td>
<td>uniform</td>
</tr>
<tr>
<td>rms bunch length</td>
<td>(\sigma_z) [cm]</td>
<td>7.55</td>
<td>7.55</td>
<td>3.78</td>
<td>20</td>
<td>6000</td>
</tr>
<tr>
<td>beta function at IP1&amp;5</td>
<td>(\beta^*) [m]</td>
<td>0.55</td>
<td>0.5</td>
<td>0.25</td>
<td>0.25</td>
<td>0.25</td>
</tr>
<tr>
<td>crossing angle</td>
<td>(\theta_c) [\mu rad]</td>
<td>285</td>
<td>315</td>
<td>445</td>
<td>430</td>
<td>1000</td>
</tr>
<tr>
<td>Piwinski parameter</td>
<td></td>
<td>0.64</td>
<td>0.75</td>
<td>0.75</td>
<td>2.8</td>
<td>2700</td>
</tr>
<tr>
<td>luminosity</td>
<td>(L) [10^{34} \text{ cm}^{-2} \text{s}^{-1}]</td>
<td>1.0</td>
<td>2.3</td>
<td>9.2</td>
<td>8.9</td>
<td>9.0</td>
</tr>
<tr>
<td>events/crossing</td>
<td></td>
<td>19</td>
<td>44</td>
<td>88</td>
<td>510</td>
<td>5 \times 10^5</td>
</tr>
<tr>
<td>rms length of luminous region</td>
<td>(\sigma_{\text{1m}}) [mm]</td>
<td>44.9</td>
<td>42.8</td>
<td>21.8</td>
<td>36.3</td>
<td>16.7</td>
</tr>
</tbody>
</table>

4.7 Example Parameter Sets

Table 2 compares the beam parameters for the nominal and ultimate LHC with those typical for three different upgrades: (1) baseline scheme with 12.5 ns spacing, (2) a scheme with large Piwinski-angle and 75-ns spacing, and (3) a superbunch. The luminosity for all upgrade schemes is close to 10^{35} cm^{-2}s^{-1}. The average beam current is 70% higher for the baseline scheme. The number of events per crossing increases from 88 for the baseline, over 510 for option (2) to an impressive 5 \times 10^5 for a single superbunch. For comparison, the nominal LHC produces about 20 events per crossing and the ultimate 45. Note also that the crossing angle is moderately increased from an about 300-\mu rad nominal value to values near 450 \mu rad in options (1) and (2) and to 1 mrad for option (3).

4.8 A Staged Approach — Review & Outlook

The LHC upgrade will likely occur in stages as outlined in the 2001 feasibility study [5]. Phase 0 of the upgrade refers to reaching the maximum performance without any hardware changes, phase 1 to the maximum performance with the arcs kept unchanged, and phase 2 to the maximum performance with ‘major’ changes. The starting point of all upgrades is the nominal performance at 7 TeV, which corresponds to a total beam-beam tune spread of \(\Delta Q_{\text{bb}}\) \approx 0.01 and to a luminosity \(L \approx 10^{34} \text{ cm}^{-2}\text{s}^{-1}\) in IP1 and IP5 (ATLAS and CMS), halo collisions or collisions with large \(\beta^*\) in IP2 (ALICE), and low-luminosity in IP8 (LHCb) \footnote{If in ALICE large-\(\beta^*\) head-on collisions occur instead of halo collisions, the beam-beam tune spread of the nominal LHC is increased and the nominal luminosity may be reduced.}. The ingredients of upgrade phases 0 and 1 correspond to the scenarios already described in previous paragraphs. They are here repeated from the perspective of staging.

The phase-0 baseline is realized by colliding the beams only in IP1&5 with alternating horizontal-vertical crossing, by increasing the bunch charge to the beam-beam limit, which is expected to be reached at \(L \approx 2.3 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}\), and, quench protection and beam losses permitting, by optionally increasing the dipole field to 9 T (ultimate field) which will raise the beam energy to 7.54 TeV.

The alternative phase-0 Piwinski scheme would increase the longitudinal emittance and rms bunch length, for example to a value \(\sigma_z \approx 15.2\) cm. The crossing angle would be enlarged by about 10%. Increasing \(N_0\) up to the beam-beam limit \((N_0 \approx 2.6 \times 10^{11})\) would allow raising the luminosity to \(3.6 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}\), without any hardware change [42].

Figure 18 compares tune footprints for the nominal LHC, the ultimate LHC (baseline phase 0) and the configuration with large Piwinski parameter (phase-0 Piwinski scheme). In all three cases, the beams are assumed to collide only at two IPs with alternating crossing. The tune spread for both upgrade variants is about the same, roughly 0.01, and, hence, comparable to the tune spread for the nominal LHC with 4 IPs. The Piwinski scheme promises about 50% more luminosity.

Phase 1 comprises various possible steps to increase the luminosity with hardware changes only in the LHC insertions and in the injector complex.

For the phase-1 baseline these steps are to modify the insertion quadrupoles and the IR layout so as to reduce \(\beta^*\) to 0.25 m, to increase the crossing angle by a factor 1.4, to increase the bunch population to the ultimate intensity, which raises the luminosity to \(3.3 \times 10^{34} \text{ cm}^{-2}\text{s}^{-1}\), to halve the bunch length \(\sigma_z\) with a higher harmonic rf sys-
One of the most difficult challenges will be to produce the high-field magnets at a reasonable cost, e.g., for less than 5 kEuros per (double) Tesla-meter, including cryogenics, which is to be compared with 4.5 kEuros per (double) Tesla-meter for the present LHC. One promising approach is the common-coil design for a double aperture magnet sketched in Fig. 19. The coils couple the two apertures and can be flat, which avoids difficult ends. In 2003, a Nb$_3$Sn dipole with a different block-coil design reached a field of 16 T, as is shown in Fig. 20.

![Figure 19: Magnet design concept for a common-coil dipole (43).](image)

Figure 19: Magnet design concept for a common-coil dipole [43].

![Figure 20: Training history of the Nb$_3$Sn dipole magnet HD-1 at LBNL (first thermal cycle) (44).](image)

Figure 20: Training history of the Nb$_3$Sn dipole magnet HD-1 at LBNL (first thermal cycle) [44].

## 5 SUMMARY AND RECOMMENDATIONS FOR FUTURE STUDIES AND R&D

Reaching the nominal LHC performance is challenging. We need to learn how to overcome electron-cloud effects, how to inject, ramp and collide almost 3000 high-intensity beams, and — if permitted by electron cloud, impedance and collimation — to double the number of bunches (and increase $\theta_c$ further) to obtain $L \approx 9.2 \times 10^{34} \text{cm}^{-2}\text{s}^{-1}$ [5].

The phase-1 Piwinski scheme would also modify the insertion quadrupoles and IR layout to provide $\beta^* = 0.25 \text{ m}$. In this scheme, the crossing angle would be increased by about 60%, and optionally every 3 bunches merged into larger bunches with 75-ns spacing. The bunch population would finally be increased to the beam-beam limit. A luminosity of $L \approx 8.9 \times 10^{34} \text{cm}^{-2}\text{s}^{-1}$ could be achieved for flat bunches with a total length $\sqrt{2\pi\sigma_z} \approx 50 \text{ cm}$ [5]. As stated earlier, a more daring option would be the generation of a 1-A superbunch per beam colliding under a large crossing angle. This, together with the smaller $\beta^* = 0.25 \text{ m}$, would provide a luminosity of $L \approx 9.0 \times 10^{34} \text{cm}^{-2}\text{s}^{-1}$.

Phase 2 of the upgrade refers to a combined luminosity and energy upgrade. To this end, the LHC injectors are modified to significantly increase the beam intensity and brilliance beyond the ultimate value, possibly in conjunction with beam-beam compensation schemes. To this end, the SPS is equipped with superconducting magnets, the transfer lines from the SPS to the LHC are upgraded in energy, and the LHC injection energy is raised to about 1 TeV. In the LHC itself, new dipoles with 15-T field and a safety margin of 2 T are installed. Such magnets are considered a reasonable target for 2015 and could be operated by 2020. The beam energy of the phase-2 upgrade would be around 12.5 TeV.

![Figure 18: Comparison of tune footprints, corresponding to betatron amplitudes from 0 to 6$\pi$, for the nominal LHC (red-dotted), ultimate (green-dashed), and large Piwinski parameter configuration (blue-solid), with alternating horizontal-vertical crossing at only two interaction points [H. Grote] [5].](image)

Figure 18: Comparison of tune footprints, corresponding to betatron amplitudes from 0 to 6$\pi$, for the nominal LHC (red-dotted), ultimate (green-dashed), and large Piwinski parameter configuration (blue-solid), with alternating horizontal-vertical crossing at only two interaction points [H. Grote] [5].
bunches, how to protect the superconducting magnets, how to safely dump the beams, and much more. Nevertheless, upgrades in beam intensity are a viable option. They require R&D for cryogenics, vacuum, rf, beam dump, and injectors, and they imply operation with larger crossing angles.

The radiation limit for the IR quadrupoles, corresponding to about 700 fb\(^{-1}\) will be reached by 2013–2014. New triplet quadrupoles with high gradient and large aperture or alternative IR layouts are needed for the luminosity upgrade. Opening the quadrupole apertures has the additional merit of intercepting less collision debris. The triplet aperture should also be large to reduce the collimator impedance for squeezed \(\beta^*\) and physics conditions.

Further studies are needed to specify the field quality of the new IR magnets, and the required upgrades of beam instrumentation, collimation and machine protection.

Experimental studies on electron cloud and on long-range or strong-strong beam-beam effects are important, and so are machine studies at existing hadron colliders with large Piwinski parameter and many (flat) bunches. An international collaboration (US-LALRP/CARE) is welcome — or needed — for LHC machine studies and commissioning.

Beam-beam compensation schemes with pulsed wires would reduce the tune footprints and the loss of dynamic aperture due to the long-range collisions. An experimental validation of such a compensation scheme is underway in the CERN SPS and may be extended to RHIC.

Interesting possibilities, currently under study, to pass each beam through separate final quadrupoles include alternate separation schemes with separation dipoles in front of the triplet quadrupoles, collision of long bunches with large crossing angle, and normal bunches at large crossing angle with crab cavities. Peak luminosities would approach \(10^{35}\) cm\(^{-2}\) s\(^{-1}\) for all schemes.

The superbunch and ‘large Piwinski angle’ options are interesting for large crossing angles, can potentially avoid electron-cloud effects, and minimize the cryogenic heat load. One could inject a bunched beam, accelerate it to 7 TeV, and then use a multiple harmonic rf system to form between 30 and 900 longer bunches. The larger the number of bunches, the smaller is the event pile up in the experiments.

Crab cavities are attractive, likely raise the beam-beam limit, and allow for separate magnet channels. First experience with crab cavities will be gained at KEKB from early 2006. Viability of crab cavities for hadron beams, e.g., the possible emittance growth due to rf phase noise, need to be further explored.

A major and sustained R&D effort on new superconducting materials and magnet design is needed for any LHC performance upgrade. For this, it is important to foster and extend collaborations with other laboratories, e.g., in the CARE and US-LARP frameworks. The development of new low-\(\beta\) quadrupoles with high gradient and larger aperture based on Nb\(_3\)Sn superconductor requires 9–10 years for short-model R&D, component development, prototyping and final production.

An increased LHC injection energy of 1 TeV in conjunction with beam-beam compensation schemes would yield an integrated luminosity gain by more than a factor of 2. A pulsed Super-SPS (and new superconducting transfer lines from the SPS to the LHC) or cheap low-field booster rings in the LHC tunnel could be the first step towards an LHC energy upgrade.

6 ACKNOWLEDGEMENT

Many CERN, US, EU and KEK colleagues have contributed to the LHC upgrade studies and the various scenarios described in this paper.

7 REFERENCES


MACHINE PROTECTION

K. H. Meß, CERN, Geneva, Switzerland

Abstract

The machine protection system of the LHC is designed to minimise the impact of beam losses and quenches. The question is discussed, whether these precautions are sufficient for the discussed upgrade scenarios.

INTRODUCTION

Machine protection consists of a complex of measures to prevent damage to the accelerator or destruction of components in case of malfunctions. As the survival of the accelerator depends on its machine protection system, these systems are usually based on a fail safe network connecting the various contributing subsystems.

The overall performance of a machine protection system depends therefore not only on the chosen network but very essentially on the installed subsystems and their performance. Some of these subsystems are discussed in their own right at this workshop separately. Hence this paper will be focussed on the overall performance of the machine protection system facing higher intensities and higher beam energies.

In protecting against damage and long repair times a machine protection system maximises inherently the beam-on time. On the other hand, it can force beam dumps accidentally. The overall machine condition, in particular beam losses, electrical noise, heat, can have an severe impact on the number of unnecessary beam dumps.

MACHINE PROTECTION OF THE LHC

The main sources of danger are the stored energy in the beam, the stored energy in the magnets and the energy coming from the mains. Consequently, the main subsystems are: the magnet powering system, the beam dump including the septa and kickers, and the extraction system for the stored magnetic energy.

The main information sources, signalling the imminent danger are a well tuned and fast beam loss monitoring system, a system to detect overheating of normal conducting magnets and cables, as well as the equivalent for the superconducting magnets, links, and current leads. The latter is usually called “quench protection system” (a euphemism par excellence as it does not protect quenches nor does it protect against quenches). In addition the information that a vital subsystem stops to function properly is also routed to the network. Finally, all other subsystems that can have an influence on the beam or the stored magnetic energy as the RF, the cryogenics, the mains power (including the UPS), the vacuum valves, the injection elements etc. contribute to the machine protection. Also beam position measurements are sometimes used as a warning signal and may be included in a future extension of the LHC machine protection system, as it is done in other machines [1]. The LHC machine protection system was recently described in ref [2]. A more explicit description of the underlying network, as it was planned at that time, can be found in ref [3].

The LHC has such intense beams that not only the experiments, but also the machine components have to be protected against the constant bombardment with lost particles and the subsequent showers. In fact, the LHC can never reach its design luminosity without an extremely and unprecedented efficient collimator system, as was pointed out by R. Assmann [4] in his talk at this workshop. In a sense the collimation system has become the first line of defence against beam losses, while the beam dump [5] serves as the last resort. This explains why we see the collimators and the beam dump system as the main issues, almost separated from the “trivial” rest of the machine protection. Consequently, the organizers of this workshop have split the whole machine protection complex in three parts: the collimators, the beam dump, and the rest. We will therefore, as far as this paper is concerned, assume that all collimator and beam dump issues are solved.

THE IMPACT OF AN INTENSITY UPGRADE

One of the steps towards a luminosity increase of the LHC consists of increasing the intensity [6]. The beam current would be raised from the design value of 0.56 A to 0.86 A or even 1.32 A. Strictly speaking, the machine protection is, under the above stated assumptions, almost not concerned. The beam loss monitoring system will have to protect against beam induced quenches by signalling significant losses early enough. In the worst case the magnets will quench and the machine protection system will dump the beam. This will of course reduce the beam-on time considerably. A quantitative estimate can not be given. The result depends entirely on the efficiency of the collimation system.

The machine protection system is such works also at increased beam intensities. There might be reliability issues for the electronics for protection of the warm magnets and a higher sensitivity to quenches of the superconducting links in the cleaning sections 3 and 7, however.

MACHINE PROTECTION FOR THE “ULTIMATE” LHC

The next step in LHC upgrade, mentioned in ref [6], involves an energy upgrade to 7.54 TeV, corresponding to
a magnetic field of 9 T in the MBs. In the reference [6] it is stated that this mode might be limited by either cryogenics or the beam dump system. As was pointed out by B. Goddard [5], the beam dump will be able to cope with this small energy increase. The higher current density and the higher magnetic field have, however, an effect on the magnet stability. Figure 1 and figure 2 (courtesy S. Russenschuch) explain this.

The load line is in the plane, which is parallel to the J-B plane. The operating point is somewhere along the line, well below the “tent”.

Temperature excursions don’t move the working point along the load line, but in a plane parallel to the J-T plane. A temperature increase pushes the working point sideways towards the “tent”, as shown in figure 2. The distance between working point and critical surface along this sideways path is in figure 3 [7] plotted as a function of the magnetic field.

![Figure 1 Load line of the LHC dipoles (Courtesy S. Russenschuch)](image1)

The following table 1 corresponds to the table 7.3 in the LHC Design Report. It shows clearly that the margin left for absorbing steady state beam losses decreases at 9 T to 1.8 mW/cm³.

![Figure 2 Temperature Margin (Courtesy S. Russenschuch)](image2)

### Table 1 Temperature margin available for steady state beam loss.

<table>
<thead>
<tr>
<th></th>
<th>@8.33T</th>
<th>@9 T</th>
</tr>
</thead>
<tbody>
<tr>
<td>Jc, precision</td>
<td>±2.5%</td>
<td>±2.5%</td>
</tr>
<tr>
<td>Dissipation</td>
<td>&lt;0.5 mW/cm³</td>
<td>&lt;0.54 mW/cm³</td>
</tr>
<tr>
<td>joints</td>
<td>0.08 0.02 K</td>
<td>0.09 0.02 K</td>
</tr>
<tr>
<td>Ramp loss</td>
<td>4.5 mW/cm³</td>
<td>1.8 mW/cm³</td>
</tr>
<tr>
<td>Beam loss</td>
<td>1.12 K</td>
<td>0.44 K</td>
</tr>
<tr>
<td>Total ΔT</td>
<td>1.4 K</td>
<td>0.8 K</td>
</tr>
</tbody>
</table>

Table 1 Temperature margin available for steady state beam loss.

The reduction of absorbable beam loss corresponds to a decrease of the number of secondaries by a factor of 2.5 and a decrease of primary lost protons by a factor of 2.7. Unless this improvement can be guaranteed by the collimation system, the time lost with quench recovery will increase dramatically.
MACHINE PROTECTION OF A 14 TEV UPGRADED LHC

This scenario is based on immature technologies. Hence a detailed assessment is not possible. However, most components of the machine protection system, as the beam loss detection system and the quench detection system, may remain unchanged. Of course, collimation and beam dump are different issues.

The most striking differences are the increase in stored energy in the beam and in the magnets.

The size of the tunnel gives an upper limit on the magnet size of 16.6 T magnets (14 TeV). The coil volume has to stay practically constant. It follows that the stored magnetic energy will go up with the square of the magnetic field by a factor 4 or slightly more.

\[ E = \frac{1}{2} I^2 \]

One can achieve this by increasing the inductance (number of windings) by a factor 4 or by increasing the current by a factor 2 or a suitable combination. Note that the magnets will still be in series and hence every magnet has to be able to absorb its own stored energy and to bypass the energy of the rest of the system between the neighbouring energy extraction resistors. This sets the requirements for the quench heaters and/or the maximum length of the magnets.

We calculate the maximum voltage at the beginning of an energy extraction to be

\[ U = I \cdot R = -L \frac{dI}{dt} \]

The current decay time constant is

\[ \tau = \frac{L}{R} \]

The maximum voltage U is most likely not far different from the present 1 kV.

Assuming a constant inductance, the current I has to go up by 2 and the dump resistance has to decrease by 2 to meet the voltage requirement. Hence, one needs in the present scheme cold diodes of 26 kA with 4 times bigger heatsinks (i.e. 240 kg). Our present experience with industry shows that industry will be reluctant to develop a cold diode suitable for 26 kA. Even 13 kA diodes might be difficult to procure. And 240 kg heatsinks are not easy to exchange in the tunnel.

Of course, one could adopt the emergency current leads (HTS?) of the Tevatron with warm diodes. Still, the energy extraction would be difficult to accommodate in the tunnel with warm switches, twice the size of the present design, and resistors four times bigger, to absorb the energy. Maybe new current HTS leads with integrated switches would fit into the tunnel.

It would be healthier to keep the maximum current and the current decay time constant by subdividing the sector not in two but in eight pieces. One would need more current leads, obviously. To accommodate the switches in the tunnel, the switches have to be part of the current leads or additional short straight sections will be needed. The energy per resistor would even be lower than presently and diodes, as are used today, could serve as bypass. Alternative powering schemes may be possible as well, but are likely to meet difficulties as well due to space limitations.

The case is less severe for the quadrupole circuits. However, a compatible design for the final focus and its protection has still to be found.

SUMMARY

The performance of the machine protection system relies on many systems and has yet to be proven. Assuming that the network part and the logic work as designed, there is little reason, why it should not work properly at higher intensities or slightly higher energies.

Collimation, monitors and beam dump have, of course, to work up to the design level. The collimation needs to surpass the performance at 7 TeV even further, to make a reasonable beam-on time possible.

In the 14 TeV case a new layout of the energy extraction system is required. This seems possible, if the magnet design keeps these boundary conditions in mind.

ACKNOWLEDGEMENT

The author thanks R. Schmidt for carefully improving the text and A. Ballarino for discussions about HTS leads and integrated switches.

REFERENCES

[7] LHC Design Report, Vol 1, p 160, Figure 7.2
VACUUM ISSUES FOR AN LHC UPGRADE

O. Gröbner*, ret. CERN, Geneva, Switzerland

Abstract
Vacuum issues for an LHC upgrade are discussed. With the decision to coat all room temperature vacuum chambers with a getter film, the vacuum performance in these sections will be significantly improved with respect to the original design. Since the activated getter film also reduces the secondary electron yield, electron cloud effects should be strongly suppressed or absent. Ionisation of the residual gas may occur not only by primary ionisation with the circulating proton beam but also by secondary ionisation with electrons of the cloud. This effect can be shown to lower the critical current for pressure stability. The slow drift of ions to the wall may furthermore create a positive space charge, which in turn can trap low energy electrons for times long compared with bunch and batch spacing in the LHC. This effect opens the possibility of a feedback between vacuum pressure and electron cloud effects.

INTRODUCTION
The LHC vacuum system is designed to meet the requirements for nominal machine parameters. After a few years of running it should be possible to achieve operation at ultimate machine conditions. For the vacuum system, this running-in period will reduce the dynamic outgassing phenomena, mainly the yields for photon, ion and electron stimulated gas desorption. Apart from electron stimulated desorption, which is the primary manifestation of the electron cloud and of beam induced multipacting (BIM), all vacuum related effects depend on the average beam current rather than on the bunched structure of the beam. From past experience, e.g. the room temperature LEP vacuum system, it is known that initial yields can be reduced by several orders of magnitude within a few months of operation [1]. For cryogenic vacuum systems the corresponding experience from laboratory systems is more limited and based on observation times equivalent to a few days of LHC operation only. Nevertheless, from the existing data one can conclude that while the room temperature photon stimulated desorption yield decreases with photon dose at least as D^{0.6} the few cryogenic measurements suggest an exponent closer to 0.3 for the LHC. Extrapolations to much larger photon doses are necessary to predict the LHC performance at ultimate conditions. Laboratory measurements of the electron stimulated desorption yield indicate a behaviour similar to the photons. Here again, data at cryogenic temperature and for significant electron doses are few or missing. The ion stimulated desorption yield (at the origin of the ISR-type pressure instability) has been measured in the laboratory at room temperature and at cryogenic temperature for some gas species. Unfortunately these yields and their evolution during the operation of the LHC are still very poorly known. For this reason, the LHC vacuum design assumes that the ion stimulated desorption remains constant during operation and the system has been designed so that the vacuum stability criterion ηI_{ultimate} < ηI_{crit} is met even under pessimistic assumptions for the ion induced desorption yield.

SYNCHROTRON RADIATION
Synchrotron radiation in the LHC can produce photon stimulated molecular desorption and photo electrons. In LHC the critical photon energy is rather low (45 eV in LHC) and all power is contained easily in the beam pipe. When the critical energy exceeds ~20 keV in some future machine, Compton scattered photons of the high energy end of the spectrum will escape and irradiate machine components (e.g. sc coils), hence shielding of the beam pipe may become an issue as for the LEP ring.

The synchrotron radiation power is

\[ P(W/m) = 1.24 \cdot 10^3 \frac{E^4(\text{TeV}) I(A)}{\rho^2(m)} \]

The linear photon flux for a proton beam is

\[ \Gamma(\text{photons/s/m}) = 7 \cdot 10^{19} \frac{E(\text{TeV}) I(A)}{\rho(m)} \]

and the photon stimulated pressure rise

\[ \Delta P \propto \eta(E_c) \frac{dN}{S} \frac{1}{dsdt} \gamma = \eta(E_c) I(A) \]

where S denotes the pumping speed per unit length (in units of m³/s). The dynamic pressure rise in the LHC, maintaining the same pumping speed and the same bending radius, depends on the molecular desorption yield \( \eta(E_c) \), beam energy and current as [2]

\[ \Delta P \propto \eta(E_c) E I \]

Desorption yield
Data on the photon energy dependence of the molecular desorption yield are scarce. A compilation of some available data is shown in Figure 1 [3, 4], giving the molecular desorption yield for H₂ and for CO as a function of the critical photon energy at room temperature and at 77 K. Room temperature molecular desorption yields appear to increase approx. proportional with the critical photon energy and hence would scale as the third
power of the beam energy. For cold measurements at 77 K, the dependence on the critical photon energy appears to be weaker, approx. as the 2/3 power and hence the yield would scale as the square of the beam energy. The scaling of the specific pressure rise is

$$\Delta P_{\text{cold}} \propto E^3 I \quad \text{and} \quad \Delta P_{\text{warm}} \propto E^4 I.$$ 

\[ \text{Figure 1: Molecular desorption yield for H}_2 \text{ and for CO as a function of the critical photon energy at room temperature and at 77 K.} \]

To improve the vacuum performance NEG surfaces have shown to be a promising solution. Commercial st707 NEG can yield an improvement of at least a factor of 10 as compared with baked stainless steel ($\eta < 10^{-5}$ molecules/photon) shown in Figure 2 [5]. According to a recent design change, all room temperature sections in LHC will have getter coatings and will be baked [6]. Coatings with a NEG-film provide a further important improvement giving $\eta \sim 10^{-6}$ molecules/photon immediately after activation and without requiring a beam conditioning time [7]. This solution combines the simultaneous advantages of low gas desorption rate & high pumping speed. Unfortunately no data are available for the dependence of the desorption yield on critical energy for the getter film.

\[ \text{Figure 2: Molecular desorption yield for commercial st707 getter as function of photon exposure as compared to baked stainless steel.} \]

**BEAM INDUCED HEAT LOAD & RADIATION DOSE**

For the LHC the power loss by nuclear scattering on the residual gas becomes a dominating effect for the vacuum requirement of the ring.

$$P(W/m) = \frac{IE}{c\tau} = 0.93 \frac{I(A)E(\text{TeV})}{\tau(h)}$$

Since the power cannot be removed by collimation, it generates in the LHC a cryogenic load for the 1.9 K cooling system. For the LHC design a limit of 0.1 W/m for two beams at ultimate current has been set equivalent to a nuclear scattering lifetime of ~100 h. The required gas density amounts to $10^{15}$ H$_2$/m$^3$ or correspondingly less, i.e. $1.2 \times 10^{14}$ /m$^3$ for CO. This loss represents a radiation dose to the s.c. magnet coils.

Laboratory measurements, e.g. with the COLDEX system in EPA at cryogenic temperature have shown that this lifetime requirement can be satisfied for the nominal beam current without any beam cleaning [8].

Using the scaling of the dynamic pressure with beam energy and with beam current, and since the beam lifetime is inversely proportional to the average gas density, the nuclear scattering losses scale as

$$\frac{dW}{ds} \propto E^a I^2$$

with $a = 4$ for the cold arcs and $a = 5$ for warm sections. To meet the requirement for the ultimate beam current, a modest improvement of the vacuum by a factor of 2.2 will be required. The scaling for the NEG coated room temperature sections is not known since desorption yield data are not available.

**ELECTRON STIMULATED DESORPTION**

Electron stimulated gas desorption by beam induced multipacting (BIM) has been observed in many machines: first in the ISR in 1977 [9] more recently in KEK-B [10], PEP-2 [11], RHIC [12] and in the SPS with LHC-type beams [13].

The gas load, $Q_{\text{cloud}}$, is related to the power deposited by the electrons, $P_{\text{lin}}$, to the molecular desorption yield, $\eta_e$, and to the average energy of the electrons in the cloud, $E_{\text{cloud}}$:

$$Q_{\text{cloud}} = k \frac{\eta_e P_{\text{lin}}}{<E_{\text{cloud}}>}.$$ 

The factor $k$ converts molecular density to pressure units. Beam conditioning, so-called scrubbing, can reduce the pressure rise by two processes: firstly by the reduction of the electron cloud power $P_{\text{lin}}$ via the reduction of the secondary electron yield of the surface and secondly by the reduction of the electron stimulated desorption yield, $\eta_e$ by surface clean-up as a function of the electron dose, see Figure 3 [14].
The parameters, which determine the vacuum stability limit for the ion induced pressure rise, are shown in Figure 4, i.e., the molecular desorption yield $\eta$ (molecules/ion), unit charge $e$ and the ionisation cross section of the residual gas by high energy protons $\sigma$. The critical current ($\eta I_{crit}$) defines the stable pressure limit.

\[ P(I) = \frac{P_0}{1 - \frac{\eta I}{e \sigma S_{eff}}} \]

$S_{eff}$ represents the effective linear pumping speed of the system. For the LHC with a beam screen the minimum pumping is provided by the pumping holes, a few % of the surface. In the room temperature sections, the NEG film assures a safety margin since it provides both, low desorption yields and large effective pumping speed. Since the relevant data for the desorption yield are missing, the improvement cannot be quantified.

**BEAM PUMPING EFFECT**

Since the bombarding ions are taken from the residual gas, it is more appropriate to introduce a net ion induced desorption

\[ \eta_{net} = (\eta - \kappa) \sigma \frac{I_b}{e} \]

where $\kappa$ represents the probability that the incident ion is captured in the wall. A negative value of the net desorption yield would imply that the capture probability is larger than the desorption and as a consequence, the beam would acts as an ion pump. Unfortunately, even for a well-designed commercial ion pump $\kappa \approx 0.1$. For a real accelerator vacuum chamber surface, the capture probability is likely to be even smaller. Beam pumping as shown in Figure 4 has been observed in the ISR as a ‘curiosity’.

In retrospect, it is important to note that in all reports from the ISR $\kappa = 1$ had been assumed. When comparing the data with measured desorption yields of baked surfaces as shown below, it remained a surprise that beam pumping could not be observed more frequently and thus this desirable effect could never be exploited systematically even with very clean vacuum chambers.

**Ion induced molecular desorption yield**

Room temperature molecular desorption yields (molecules/ion) as a function of the ion energy for unbaked and baked stainless steel are shown in Figs. 5 and 6. A laboratory program is ongoing to obtain the relevant data for cold surfaces [16].
In a vacuum system the residual gas is a mixture of different species and it is therefore important to know the mutual desorption yields of molecules. From past experience of room temperature vacuum systems, the species which dominate the ion induced pressure rise have been CO rather than H$_2$ because of the larger ionisation cross section and the smaller pumping speed. For the design of vacuum systems, which are mainly pumped with getter pumps, other species like CH$_4$ or Ar, which are not chemically pumped, could become important.

**VACUUM STABILITY IN A CRYOGENIC VACUUM SYSTEM**

Molecules are cryo-pumped with high efficiency directly onto the cold bore. The pumping speed per unit length is

\[ S_{\text{eff}} = \frac{1}{4} \nu s F \]

with \( \nu \) the mean molecular velocity, \( s \) the sticking probability of molecules on the wall and \( F \) the surface area per unit length. With the LHC beam pipe radius \( r_p \), the stability limit is

\[ (\eta)_{\text{crit}} = \frac{\pi}{2} \nu s r_p \frac{e}{\sigma} \]

and for \( s \sim 1 \) the critical current is of the order of kA. However, this large stability limit is offset by two factors: firstly, the sticking probability is less than unity and secondly, the molecular desorption yield, \( \eta \) for thick layers of condensed gas, specifically for H$_2$, which accumulates on a cold wall can become very large, up to \( 10^4 \) molecules per ion as shown in Figure 8.

**IONISATION OF THE RESIDUAL GAS**

*Primary ionisation of the residual gas*

The ionisation cross section of CO for high energy protons is \( \sigma_p \sim 2 \times 10^{-22} \) m$^2$. The circulating beam current produces an ion current to the wall (A m$^{-2}$)

**H$_2$ vapour pressure increase by exposure to thermal or to synchrotron radiation**

In the course of the initial studies for the SSC vacuum system [17] and later during the measurements for the LHC [18] it became clear that cryo-sorbed molecules, most specifically H$_2$ are re-desorbed readily by low energy synchrotron radiation and even by thermal radiation. This effect will make it impossible to pump H$_2$ in the required quantities on a ‘naked’ cold bore exposed to synchrotron radiation. This limitation of cryo-pumps due to the exposure to environmental room temperature radiation and to the bombardment of beam induced energetic particles (photons, electrons, ions) must be taken into account in any design. For commercial cryo-pumps this effect imposes the installation of LN$_2$ cooled baffles and for the LHC it imposes a beam screen. It should be noted that this requirement arises not only for heat load reasons but mainly to avoid re-desorption of weakly adsorbed molecules. It merits to be noted that the unexpected pressure increase by thermal radiation has been investigated during the 1970 for the cryo-pumps in the ISR and that the findings had been published [19].
Primary ionisation has been responsible for the ion induced pressure rise effect in the ISR. For the cold LHC, assuming a total CO pressure of about $10^{-7}$ Pa, one finds $I_p < 50$ nA/m. Hence it represents a small current and an insignificant power load. For LHC-type beams H$_2$ ions gain ~220 eV and CO ions ~ 180 eV. Due to the repulsion from the positive space charge of the proton beam, the relatively heavy CO ions take ~ 1µs to reach the wall, i.e. the beam screen. The light ions like H$_2$ are faster. As a consequence ions have a 'long' memory as compared to the typical bunch and batch spacing in the LHC.

Secondary ionisation of the residual gas

If the pressure increases due to the electron cloud and BIM, electrons of the cloud may in turn ionise the residual gas and the possibility of a feedback mechanism, similar to the conventional ion induced pressure bump can occur. This secondary effect represents an additional source of ions and electrons as well as a heat load for the beam screen. Furthermore, even a relatively small positive space charge can trap low energy electrons. Because of the pressure dependence, secondary ionisation will be enhanced by a beam induced pressure rise initially triggered by electrons or ions.

The secondary ionisation current by the electron cloud can be calculated as the ion current in a conventional ionisation vacuum gauge

$$I^+ = s_i L_e P_g I_e$$

Here $I^+ (A m^{-1})$ is the ion current to the wall, $s_i (m^{-1} Pa^{-1})$ is the specific ionisation of the residual gas, $L_e (m)$ the path length of electrons, $P_g (Pa)$ the residual pressure and $I_e (A m^{-1})$ the electron current of the cloud. One finds that most electrons have near optimum energy (between 50 – 300 eV) for gas ionisation with $s_i$ values of ~3 (m$^{-1}$ Pa$^{-1}$) for CO and ~ 1 (m$^{-1}$ Pa$^{-1}$) for H$_2$ [20]. The corresponding cross sections are $\sigma_S \sim 1.2 \times 10^{-20}$ m$^2$ for CO and ~ 4.0 $\times 10^{-21}$ m$^2$ for H$_2$. The unknown path length of the electrons is obtained in the case of ionisation gauges by a calibration with a known pressure.

Unfortunately, during the recent beam studies in the SPS no suitable current monitors have been available to measure this ion signal in situ since these monitors could not differentiate between ions and secondary electrons.

ION INDUCED PRESSURE RISE WITH ELECTRON CLOUD

The gas load per unit length of vacuum system i.e., $P S_{\text{eff}} (Pa m^2 s^{-1} m^{-1})$ is the sum of:

- Thermal outgassing $P S_{\text{o eff}}$
- electron stimulated desorption $\eta_e kT \frac{I_e}{e}$, primary and secondary ion desorption

\[
\eta_i \left( \frac{I_B}{p} + \frac{\sigma}{s} L_e \frac{I_e}{e} \right) P
\]

This gives

\[
P(\eta_B) = \frac{P S_{\text{eff}} + \eta_e kT \frac{I_e}{e}}{kT}
\]

Equilibrium gas density as a function of beam current can be expressed as

\[
P(\eta_B) = \frac{P S_{\text{eff}} + \eta_e kT \frac{I_e}{e}}{kT}
\]

Secondary ionisation by an e-cloud can reduce the critical beam current and hence the vacuum stability. It is interesting to make an order of magnitude estimate for this effect: -the ionisation cross section for low energy electrons is about two orders of magnitude larger than for high energy protons. The electron current during multipacting studies in the SPS with an LHC–type beam has been of the order of 0.05A/m [21] and the nominal LHC beam will be 0.54A. Hence, to make the contributions of primary and of secondary ionisation equal, that is to half the critical current, the path length of electrons would have to be 0.1m only, i.e. roughly the diameter of the circular SPS beam pipe. If electrons would have a longer path length, the secondary ionisation term could easily become the dominant effect.

Finally, since the secondary ionisation rate is proportional to pressure, BIM may initiate a vacuum instability. On the other hand, electron stimulated desorption, which is described by the second term in the numerator, produces a pressure increase, but does not lead to a feedback process and to the ion induced pressure instability.

ION SPACE CHARGE

Since ions take of the order of $10^6$ s to reach the wall, this represents a positive space charge in the form of an ion cloud filling the beam pipe. Furthermore, since the ionisation rate increases with the pressure, such an ion cloud would grow with the pressure.

$$U = \frac{e}{4 \pi \epsilon_0} \rho_i$$

One would expect that the ion cloud could in turn trap a significant number of low energy electrons. In fact, a density of a few $10^9$ ions/m$^3$ would be sufficient to trap the very low energy secondary electrons during a time scale of µs as given by the slow ions. Depending on the density, the ion cloud may increase the scattering of the beam. Finally, since this secondary ionisation effect is proportional to pressure, once a pressure rise has been
initiated, it could stay for seconds, depending the pump down speed of the vacuum system.

ION IMPACT ENERGY AT THE COLLISION POINT

Due to the strong beam focusing at the crossing points of the LHC the ion energy is increased from a few hundred eV to several keV. In addition, in the presence of a magnetic solenoid field of a detector magnet, ions must spiral to the wall and may gain an energy of

$$E_{ion} = \frac{e}{2m} B^2 r^2_{pipe}.$$  

Ion impact energies may become rather large in particular with grazing incidence and thus could induce sputtering of wall material. Instead of a conventional pressure bump one could generate a ‘metal cloud’ or a ‘metal curtain’ depending on the sputtering yield.

OUTLOOK AND CONCLUSIONS

Dynamic vacuum issues should remain a priority for an LHC luminosity upgrade. Several design parameters are not sufficiently known e.g. ion induced desorption yields as a function of ion species, energy and temperature. Getter coated beam pipes in warm sections of LHC provide a comfortable safety margin against desorption phenomena caused by electrons, photons and ions but detailed data are missing. The possibility that secondary ionisation of the residual gas can be linked to BIM and electron cloud effects should be further studied. The interaction between the space charge of ions and electrons may complete the understanding of some observations of the electron cloud. The contribution of ions should be included in the heat load budget and in this respect calorimetric measurements would be more appropriate as compared to current measurements on strip electrodes at the wall of the beam pipe. In this respect, a dedicated ion current monitor with adequate sensitivity should be designed and installed in the LHC.

ACKNOWLEDGEMENTS

I would like to thank Drs V. Baglin, M. Jimenez, J.-M. Laurent, N. Hilleret, P. Strubin and F. Zimmermann for stimulating and fruitful discussions which have guided me in the preparation of this paper.

REFERENCES

[1] O. Gröbner, Vacuum, 43 (1/2), 1992
[16] N. Hilleret, private communication, 2004
[18] V.V. Anashin, R. Calder, O. Gröbner, O.B. Malyshov, VACUUM 53 p.269, 1999
[21] V. Baglin, private communication, 2004
Space Charge and Optics Studies for High Intensity Machine

G. Franchetti *, I. Hofmann *, M. Giovannozzi †, M. Martini†, E. Métral†
* GSI, 64291 Darmstadt, Germany
† CERN, 1211 Geneva 23, Switzerland

Abstract

In the FAIR [1] facility planned at GSI high space charge may play an important role for limiting the foreseen nominal machine performance. In the SIS100, bunches with tune shift as large as ≈ 0.2 should be stored for hundred thousand turns keeping the loss budget within few percent to avoid quenching of superconducting magnets and vacuum degradation. Usually, design criteria to meet these beam dynamics constraints rely on quantities such as dynamic aperture or nonlinear acceptance. The presence of space charge in a nonlinear lattice sets, however, new challenges to beam loss control and requires an effort to understand unexplored leading beam loss mechanisms. In this context beam loss studies for a high intensity bunched beam have been performed at the CERN-Proton Synchrotron. We discuss the experimental results and propose a mechanism for explaining the observed beam loss. We then apply to the SIS100 what we have learnt from CERN-Proton Synchrotron measurements and discuss its consequences.

INTRODUCTION

We are reviewing in this proceeding the work published in the 33rd ICFA workshop and in EPAC 2004 on space charge and lattice nonlinearities beam loss. The trapping of particles into resonance-induced phase space islands has first been considered for single passage through a resonance. The efficiency of the process has been investigated in [2, 3] and mechanisms for detrapping have been studied as well [4]. Applications of the trapping of particles have been proposed as a method to clean the beam from halo particles [5] and recently for multi-turn extraction [6, 7]. Contrary to a "useful" controlled trapping/detrapping of particles, in the beam loss issue for high intensity we are dealing with uncontrolled trapping/detrapping phenomena. In the SIS100 [8, 9], for example, storage for 1 second of a high intensity bunched beam (ΔQx ≈ 0.2) in a nonlinear lattice with a loss level not exceeding 1% is requested. The proposed mechanism is as follows: when the tune Qz0 is set close above a single nonlinear resonance stable islands appears at a certain amplitude corresponding to the space charge detuning. Contrary to the single passage through a resonance, in which the island control is achieved by appropriately changing the machine tune, in the high intensity bunch space charge is now controlling the island position along the longitudinal axes. Islands have to be further out in the phase space in the center of the bunch (z = 0) to compensate the stronger space charge detuning, whereas in the head/tail (z = ±z_{max}) the island will be further in due to the weaker space charge. Hence some particles will be periodically crossed by the lattice induced space charge controlled islands. These particles then perform a periodic resonance crossing at 4 times the synchrotron frequency. In fact, a particle crosses 2 times the center of the bunch (z = 0) per synchrotron oscillation, therefore if the resonance condition is met, for instance, at quarter of the bunch length, then the resonance condition is crossed 4 times the synchrotron oscillation with the associated possible trapping-detrapping phenomena. The complexity of the single particle dynamics in the self-consistent field is considerable and so is computer simulation. A study of trapping conditions has been presented in [10]. A simplified approach was using a model with a frozen coasting beam with a Gaussian distribution and a parametrically modulated intensity [11]. Unfortunately fully self-consistent simulations are affected by the artificial noise created by the PIC solvers [12]: long term simulations (thousands of machine turns) may transfer the usually acceptable PIC noise into unphysical emittance growth or excessive halo formation. Therefore we have proposed a long term simulation using analytical expression for the electric field for a frozen charge distribution presented in detail in [13] and applied to SIS100 in Ref. [14]. In order to check the validity of these simulations a campaign of experiments has started in 2002 at the CERN-Proton Synchrotron (PS) in a CERN-GSI collaboration. The aim of these measurements was two-fold. One goal is to investigate the new mechanism of beam loss/emittance increase and provide experimental data for code comparison. The second is the investigation of the Montague resonance with a purpose of a code benchmarking [15, 16]. The first code comparison presented in [13] employed a constant focusing lattice and was only able to confirm the experimentally observed emittance growth, but not the measured beam loss. In the present paper we extend the simulation to the full AG focusing lattice of the CERN-Proton Synchrotron, in order to get a better match of experiment and simulation condition we are also referring to recent experimental data obtained in 2003.
The measurements were all performed keeping the kinetic energy at the injection value of 1.4 GeV. Bunch profiles were first measured at 180 ms after injection (before powering the octupoles). Profiles were found to be Gaussian in both vertical and horizontal. The chromaticity was close to the natural one, and the effect of the momentum spread of $3.3 \times 10^{-3}$ (at $2\sigma$) was found to be 27% of the maximum space charge induced tunespread $\Delta Q_x = 0.075$ for a particle with small amplitude. At 180 ms after injection two calibrated octupoles each with integrated strength of $K_3 = 0.6075 \times I \text{ m}^{-3}$, with $I$ octupole current, were powered to -20 A in order to excite the resonance $4Q_{x0} = 25$. The measurement window was set to 1.2 s ($4.4 \times 10^7$ turns) during which the bunch intensity was monitored with a current transformer. At selected times we measured transverse beam profiles with the flying wire (20 m/s), fitted them with a Gaussian profile, and determined the corresponding rms emittances. In most cases profiles were actually found quite close to Gaussian. The vertical machine tune was set to $Q_{y0} = 6.12$, and the horizontal tune was varied in the interval $6.24 < Q_{x0} < 6.32$ so as to link the octupole resonance crossing with different positions of the space charge tune-spread. The parameters of the experiment are summarized in Table 1. In Fig. 1 we plot, as a function of $Q_{x0}$, the emittances and beam intensity 1.2 s after injection. In the interval $6.28 < Q_{x0} < 6.32$ we find the typical emittance growth regime, which is characterized by the maximum emittance growth of 42% at the tune $Q_{x0} = 6.28$. The beam loss regime, instead, is located in the interval $6.25 < Q_{x0} < 6.28$ and the maximum beam loss of $\approx 32\%$ is found at $Q_{x0} = 6.265$. We also studied the time dependence of the longitudinal bunch profile. These results shown here in a standard waterfall plot (Fig. 2a) suggest that the lost particles make the longitudinal distribution shrinks in amplitude and size. In order to quantify this visual pattern we computed the time evolution of the bunch length by doing a Gaussian fit for each of the profiles of Fig. 2a. These results are plotted in Fig. 2b. In the same picture we plot also the integrated intensity of each longitudinal bunch profile for each time measurement. The two curves in Fig. 2b show that there is a direct relation between beam loss and bunch shortening. Unless an unexpected transverse-longitudinal correlation takes place, Figs. 2 a,b suggest that lost particles are not only those with large transverse amplitude, but also with large synchrotron amplitude. This experimental evidence is consistent with the condition of periodic resonance crossing. In fact particles which posses small longitudinal amplitude and small transverse amplitude will always have an effective tune $Q_x$ sufficiently below 6.25 because their motion is confined in the denser region of the bunch, hence they will never be extracted. Particles with large longitudinal amplitude, instead, can periodically cross the resonance and therefore may get trapped and eventually get lost.

### Table 1: Summary of experimental settings

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
<th>Units</th>
</tr>
</thead>
<tbody>
<tr>
<td>Kinetic energy</td>
<td>1.4</td>
<td>GeV</td>
</tr>
<tr>
<td>Particles per bunch</td>
<td>$1 \times 10^{12}$</td>
<td></td>
</tr>
<tr>
<td>Bunch length ($4\sigma$)</td>
<td>180</td>
<td>ns</td>
</tr>
<tr>
<td>Emittances (norm. at $2\sigma$)</td>
<td>$25.5/10$</td>
<td>mm mrad</td>
</tr>
<tr>
<td>Momentum spread ($2\sigma$)</td>
<td>$3.3 \times 10^{-3}$</td>
<td></td>
</tr>
<tr>
<td>Derived maximum tuneshifts</td>
<td>$0.075/0.12$</td>
<td></td>
</tr>
<tr>
<td>PS circumference</td>
<td>628</td>
<td>m</td>
</tr>
<tr>
<td>Beam pipe axes</td>
<td>14/7</td>
<td>cm</td>
</tr>
<tr>
<td>PS superperiodicity</td>
<td>10</td>
<td></td>
</tr>
<tr>
<td>Octupole current</td>
<td>-20</td>
<td>A</td>
</tr>
</tbody>
</table>

### Simulation

3D simulations so far have been performed by using a simplified axisymmetric frozen bunch model with density profile $\rho \propto (1 - t)^2$ [17]. Here $t = (r/R)^2 + (z/Z)^2$, with $r = \sqrt{x^2 + y^2}$, $R$ the transverse beam edge radius and $Z$ the longitudinal edge. This analytic frozen bunch model has the advantage that noise free space charge forces are obtained. In the simplified axisymmetric model the horizontal and vertical axes were artificially equal, therefore we have progressed to modeling ellipsoidal bunches [18]. For the simulations reported here we used an ellipsoidal Gaussian distribution $\rho \propto \exp(-t/2)$, with $t = (x/\sigma_x)^2 + (y/\sigma_y)^2 + (z/\sigma_z)^2$, and $\sigma_x, \sigma_y, \sigma_z$ arbitrary rms bunch sizes. The tracking is performed by computing the average beam size in horizontal/vertical axes, then the space charge algorithm is initialized keeping the 3 bunch

![Figure 1: Experimental findings: normalized final emittance and beam intensity as function of the working point.](image-url)
axes frozen. This modeling is neglecting the local transverse modulation of the bunch envelopes which we expect to have only a minor influence on the island position due to its rapid oscillation. A further improvement in simulations has been reached by using the AG focusing PS lattice as provided in [19], including the lattice nonlinearities as discussed in [20]. With this modeling the beam is now experiencing the correct nonlinear forces from the octupole.

**Dynamic aperture and space charge** In previous work [13] the PS lattice was simulated in a constant focusing approximation. This simplified the particle dynamics, but it also introduced enhanced horizontal and reduced vertical octupolar nonlinear forces at the location of the external octupoles as in this case $\beta_x = \beta_y = 16 \text{ m}$. In contrast, by using the real PS lattice, the beta functions are $\beta_x = 11.7 \text{ m}$, $\beta_y = 22.2 \text{ m}$ at the location of the octupoles. It becomes then necessary firstly to investigate if beam loss may be attributed to an excessive shrinking of the dynamic aperture (DA). To explore roughly the DA of the PS we have carried out a numerical test by searching the maximum stable radius of test particles placed into 20 different directions in the upper half of the $x-y$ plane. The border of stability in each direction is found by using a bisecting method. The iterative process stops when the uncertainty over the radius is $1\%$. In all these calculations space charge is not included. The stability condition was searched within $10^4$ turns (short term DA). We firstly computed the DA without external octupole ($I = 0 \text{ A}$), and found that it is at $\sim 9 \sigma$, for all the tune range, where $\sigma$ is the horizontal rms beam size of the injected beam. This value of the DA sets the inner most part of the stable region practically on the beam pipe, i.e. the mechanical acceptance is all inside the stable region. By activating the octupoles with $I = -20 \text{ A}$, the short term DA is lowered at $\sim 8 \sigma$ for $6.27 < Q_{x0} < 6.32$, but near $Q_{x0} = 6.25$ it shrinks to $3.5 \sigma$ (Fig. 3). This value is small enough to intercept the tail of a Gaussian distribution. However, a 2D multiparticle simulation (without space charge) on the same time scale has shown no particle loss. An upper bound to the beam loss can be obtained by cutting the 2D Gaussian distribution in energy at $3.5 \sigma$, i.e. keeping only particles such that $\epsilon_{x,s}/\epsilon_{x} + \epsilon_{y,s}/\epsilon_{y} < 3.5$, with $\epsilon_{x,y,s}$ the single particle emittances, and $\epsilon_{x/y}$ beam rms emittances. This estimate gives only $0.5\%$ beam loss. A further inspection at $10^5$ turns shows that close to $Q_{x0} = 6.25$ the DA without space charge shrinks from $3.5 \sigma$ to $3 \sigma$. Again these results suggest that the DA alone as modeled here and ignoring space charge cannot explain beam loss. We also note the significant difference between the study of the single particle stability, where the minimum DA is located at $Q_{x0} = 6.25$, from the experimental finding of Fig. 1, where the maximum loss occurs at $Q_{x0} = 6.265$. In order to understand the origin of this difference we studied a simplified constant focusing 2D model of the PS with the nonlinearities of one magnet. In this model the longitu-

Figure 2: Measurements at $Q_{x0} = 6.265, Q_{y0} = 6.12$: a) Waterfall plot of the longitudinal profile as function of time; b) Beam intensity and bunch length as function of storage time. The solid lines are fitted with data by using the functional dependance reported.

Figure 3: Dynamic aperture as function of $Q_{x0}$ with and without octupole.
dinal motion is kept frozen. The aim of this model is to study the nonlinear tune of a test particle as function of its initial coordinates \( x = x_{\text{max}}, z = z_{\text{max}}, x' = z' = 0 \).

The tune was computed with an FFT in 4096 turns [21]. In Fig. 4a we computed this "frequency map" for the bare tune of \( Q_{x0} = 6.265 \). We see that for \( x \sim z \sim 0 \) the nonlinear tune \( Q_x \) has the lower value \( Q_x = Q_{x0} - \Delta Q_x \) with \( \Delta Q_x = 0.075 \). At the longitudinal position of \( z = 3 \sigma_z \) the space charge is quite reduced and only the detuning due to the octupole is present. Note that detuning of octupole and space charge act in the same direction (they reduce the nonlinear tune) but in a different way: the space charge detuning is stronger off axis whereas the octupole detuning is stronger on axis as it can be seen at \( z = 3 \sigma_z \). The flat-top in the picture represents all the initial conditions of particles whose tune is locked on the 4th order resonance. It is approximately given where the space charge detuning matches the octupole induced resonance condition. It is interesting to note that this flat-top does not reach the plane \( z = 0 \), but stops at \( z = 1 \) where a chaotic region is met. We also note that the flat-top reaches the maximum extension of \( 5 \sigma \) far below the DA of \( 10 \sigma \). However the chaotic region extends from \( 5 \sigma \) to the DA. This picture suggests that particles trapped follow the flat-top, and when they reach the chaotic region they can be brought to the DA and get lost. This case occurs only at the tune \( Q_{x0} = 6.265 \). For comparison we have calculated the same "frequency map" for \( Q_{x0} = 6.26875 \) (Fig. 4b). In this case the flat-top crosses the plane \( z = 0 \) at \( 5 \sigma \). This becomes consequently the maximum extension of trapped particles which cannot be lost.

### 3D simulations with synchrotron motion

Using the AG structure of the PS ring increases considerably the CPU time required for simulations. For this reason we have restricted the number of turns to \( 1.5 \times 10^5 \). Following the procedure of the experiment we have performed simulations for the relevant working points used in the measurements. The initial bunch distribution was Gaussian with 2000 macroparticles with transverse emittances \( \epsilon_x = 25.5 \) mm-mrad, \( \epsilon_y = 10 \) mm-mrad normalized at \( 2 \sigma \). Chromatic effects have been neglected here. In the simulations, the PS beam pipe (14/7 cm) has been assumed constant throughout the ring. In Fig. 5 we show the result of the simulations at \( 1.5 \times 10^5 \) turns and plot for convenience also the final experimental data. The comparison shows a similar pattern: an emittance growth for larger tunes next to a loss dominated regime closer to \( Q_x = 6.25 \). In the emittance growth regime the agreement between measured and simulated data, taken at the same number of turns, is reasonably good. This is shown in Fig. 6 for the working point corresponding to the maximum emittance growth in the experiment. The discrepancy is larger in the loss dominated regime, where the simulation gives a maximum 8% for \( Q_{x0} = 6.26 \), while the maximum measured beam loss is found to be \( \approx 25 \% \) at \( Q_{x0} = 6.265 \) both after \( 1.5 \times 10^5 \) turns. This suggest that we should re-examine all approximations made in simulations, but also the accuracy of the modeling of the nonlinear lattice for the working points here used. The effect of chromaticity must be considered as well since it enters directly into the resonance condition and therefore into the trapping/detrapping probability as well as into the maximum halo radius. It should also be noted that experimental beam loss of the order of tens of percent leads this study into a realm, where self consistent simulation would be needed. To further support the interpretation of our findings we have studied the correlation between beam loss and bunch shortening, which is shown in Fig. 7. The bunch shortening is comparable with the experimental results of Fig. 2b at \( 1.5 \times 10^5 \) turns, even if for a slightly different bare tune (in fact the experimental beam loss data are relatively flat between \( Q_{x0} = 6.26 \) and \( Q_{x0} = 6.265 \)). Note that although the measured and simulated emittance and bunch length follow a similar evolution in the first 0.5 seconds, the beam loss is quite different.
APPLICATION TO SIS100

SIS100 Nonlinear lattice

This section describes our progress in DA calculations for the SIS100 using a preliminary lattice and working point [8]. In the SIS100 the storage of a beam is foreseen, which fills a good fraction of the elliptical beam pipe of semi-axes $55 \times 25$ mm [8]. Multipolar expansion of the fields generated by accelerator magnets is used to model lattice nonlinear components. Usually multipoles $b_n, a_n$ are retrieved by integrating the magnetic field over a reference circle of radius $R$. The standard calculation of $b_n, a_n$ made with $R \leq 25$ mm leaves uncertain the accuracy of the reconstructed magnetic field in the far region of the magnet aperture. Alternatively, the multipoles can be obtained through a fit [22] of the magnetic field map [23] using a truncated multipolar expansion. The best fitting provides the order of the expansion. Table 2 shows the multipoles for an SIS100 dipole at 10% excitation computed with standard [23] and fitting procedure (here $a_n = 0$). In order to improve the convergence between the standard and fit $b_3$, the magnetic field map has been taken in $|x| \leq 50$ mm. The maximum reconstruction error in the grid magnetic field is $\Delta B/B_0 = 0.38 \times 10^{-4}$. We use the fit multipoles to describe the nonlinear components in the center of the SIS100 bends. The effect of the sagitta of 8 mm [24] is included through a proper coordinate shift at the location of the bend nonlinearities. The nonlinear components of the quadrupole fringe field are modeled as well [25]. The nonlinearities of the body of the quadrupoles are not included yet.

Stability without space charge

In Fig. 8a) we plot a cut of the stability domain in the x-y plane. In this calculation we have neglected off-momentum effects. The SIS100 is tuned on the preliminary reference working point [8] of $Q_{x0} = 15.9, Q_{y0} = 15.7$. The stability domain is computed in the x-y section with $\beta_x = 7$ m
Table 2: Standard and best fit multipoles in units of $10^{-4}$.

<table>
<thead>
<tr>
<th>n</th>
<th>$b_n$(standard)</th>
<th>$b_n$(fit)</th>
<th>deviation %</th>
</tr>
</thead>
<tbody>
<tr>
<td>3</td>
<td>1.56</td>
<td>1.65</td>
<td>-5.1</td>
</tr>
<tr>
<td>5</td>
<td>$-1.11 \times 10^{-1}$</td>
<td>$-1.41 \times 10^{-1}$</td>
<td>-21.4</td>
</tr>
<tr>
<td>7</td>
<td>$1.40 \times 10^{-2}$</td>
<td>$2.43 \times 10^{-2}$</td>
<td>-42.5</td>
</tr>
<tr>
<td>9</td>
<td>$6.0 \times 10^{-3}$</td>
<td>$6.67 \times 10^{-4}$</td>
<td>800</td>
</tr>
<tr>
<td>11</td>
<td>$-4.0 \times 10^{-3}$</td>
<td>$7.7 \times 10^{-4}$</td>
<td>-617</td>
</tr>
<tr>
<td>13</td>
<td>$-4.0 \times 10^{-3}$</td>
<td>$1.75 \times 10^{-4}$</td>
<td>2184</td>
</tr>
</tbody>
</table>

and $\beta_y = 9.3$ m (injection). Each point in the plot represents the initial condition of test particles with $x' = y' = 0$ which are stable (bounded) during the turns correspondent to its grey scale. In red is plotted the beam pipe in the SIS100 bends. In blue (inner ellipse) is drawn the space section of the beam at $2\sigma$ if injected with equilibrated emittances of $\epsilon_{x,\text{rms}} = \epsilon_{y,\text{rms}} = 8.75$ mm mrad [8]. The DA is defined in Ref. [26, 27, 28] as the radius whose area is equal to the area of the stability domain. Since in SIS100 the beam is planned to fill the pipe aperture we consider the DA as the radius (in normalized coordinates) of the largest circle inscribed inside the domain of stable initial conditions [29]. As customary we express the DA in terms of the beam $\sigma$, which for equal emittances [8] reads $4.4\sigma$. This result, however, does not account for the imperfections which affect the systematic strength of each magnet nonlinear component. We assume these perturbations are of the same kind for each magnet: with average zero and variance 10% [30] of the systematic strength $b_n$. By giving each nonlinear component of the SIS100 its perturbed strength, we have formed an error set. For each error set a new DA will be found. We computed the DA for $5 \times 10^4$ turns and 97 error sets. These results are shown in the histogram in Fig. 8b). The vertical axes gives the number of sets for which the DA aperture was found to have the value in the correspondent bin in the horizontal axes. Our results for the DA cannot be taken as final values since additional factors need to be taken into account. For a preliminary discussion we assume the suggested reductions (in %) from Ref. [29] and scale in Table 3 the DA-values at $4.4\sigma$ accordingly (for the reference working point). We also infer possible beam loss by simply using a Gaussian-cut model. The additional safety margin of $-20\%$ proposed in Ref. [29] appears to be unacceptable here and needs further consideration.

Table 3: Assumed further reductions of computed DA and accumulated loss.

<table>
<thead>
<tr>
<th>Source of Uncertainty</th>
<th>Suggestion</th>
<th>DA</th>
<th>% Loss</th>
</tr>
</thead>
<tbody>
<tr>
<td>Sensitivity to initial condition</td>
<td>-10%</td>
<td>3.96</td>
<td>0.3</td>
</tr>
<tr>
<td>Time-dependent multipole</td>
<td>-10%</td>
<td>3.65</td>
<td>1.3</td>
</tr>
<tr>
<td>Ripple</td>
<td>-10%</td>
<td>3.28</td>
<td>3.6</td>
</tr>
</tbody>
</table>

Figure 8: a) Stability domain; b) DA for a beam with $\epsilon_{\text{rms,x,y}} = 8.75$ mm mrad as function of bend errors; c) DA versus working point.

Effect of space charge on long term bunch storage

In the FAIR project [8] it is foreseen that a primary beam $U^{28+}$ is injected into the SIS100 from the upgraded version of SIS18 at 96 MeV/u. There, an RF system with harmonic
10 is activated and 8 of the buckets are filled. The total ions stored in the SIS100 is planned to be $10^{12}$ in 1 second. This storing time corresponds approximately to $10^5$ turns. We consider here a reference bunch with unnormalized transverse emittances at $2\sigma$ of $\epsilon_x = 50$, and $\epsilon_y = 20$ mm-mrad containing $1.25 \times 10^{11}$ uranium ions. The momentum spread at $1\sigma$ is $2.5 \times 10^{-4}$. We consider here a typical frequency of 1000 turns per synchrotron oscillation. In this scenario a high intensity bunch is stored for long time and special attention must be paid to the interplay between lattice nonlinearities and nonlinear space charge forces. As discussed in Section 1, synchrotron motion may reduce beam quality as it induces trapping/detrapping phenomena. Beam emittance increase and particle loss are the final consequences of such a mechanism. The most dangerous effects happens when the bare tune is just above a resonance line as in this case synchrotron motion may induce - through space charge - a periodic crossing of this resonance at large transverse amplitude. Of critical importance for emittance growth and beam loss are the two quantities: 1) the distance of the bare tune from an excited resonance $\Delta Q_\pi$; 2) the maximum tuneshift $\Delta Q$. Extraction of particles may happen when $\Delta Q/\Delta Q_\pi$ is sufficiently large. Crucial is also the number of synchrotron oscillations performed during the storage time which has been explored up to 100 synchrotron oscillations. A first inspection of Fig. 8c reveals that the SIS100 reference working point $Q_{x0} = 15.9$, $Q_{y0} = 15.7$ is located above the systematic 5th order resonances $Q_x + 4Q_y = 78$, $3Q_x + 2Q_y = 78$. In spite of the systematic nature of these resonances one would expect a weak effect due to trapping and detrapping because of the high order of such resonances. It should however be pointed out that in Section 1 only one resonance was strongly excited and that here we are in presence of combined nonlinearities of several orders. The overall dynamics is therefore more complex. For a preliminary study of the effect of the synchrotron and space charge induced beam loss we consider two types of bunches.

A) A bunch with moderate space charge, subjected to a tuneshift $\Delta Q_y = -0.1$. This tuneshift is consistent with a two harmonics RF system which flattens the bunch profile. For such a tuneshift Fig. 8c shows that apparently no resonance of fifth order is crossed for the above working point.

B) A second type of bunch with stronger tuneshift is considered. For larger tunespread as $\Delta Q_y = -0.26$ the possibility of synchrotron induced trapping becomes more serious as $\Delta Q_y/\Delta Q_\pi \simeq 2$. The tunespread would cross the resonance $Q_x + 4Q_y = 78$.

The space charge effects are introduced in the tracking through an analytic frozen model of the bunch [18]. The analytic modeling prevents artificial noise effect in the tracking. For these simulations we set 50 space charge kicks per lattice period. The beam pipe is an ellipse of $55 \times 25$ mm kept constant along the ring. We consider the SIS100 tuned at the reference working point and tracked 2000 test particles in a bunch for $10^5$ turns. In order to model the moderate space charge case we consider an artificially enhanced bunch length of 130 m ($1\sigma$), which yields in the bunch center the tuneshifts of $\Delta Q_x = -0.068$, $\Delta Q_y = -0.1$. The strong space charge case is created with a bunch length at $1\sigma$ of 54 m, hence $\Delta Q_x = -0.16$, $\Delta Q_y = -0.26$. For the moderate space charge bunch

![Figure 9: Moderate space charge case results: a) evolution of transverse rms emittances and b) % of initial particles remaining in the bunch.](image-url)
tion of horizontal emittance as major loss of particles in the horizontal plane. This result seems however to contradict the initial decrease found in Fig. 9, as there the bigger emittance reduction happens in the vertical plane. We may question if this is due to pure lattice nonlinearity or from the combined effect of space charge defocusing jointly with the lattice nonlinearities.

For the same simulation but now including only lattice nonlinearities we find an initial beam loss of 8.6% and \( \epsilon_x/\epsilon_{x0} = 0.77, \epsilon_y/\epsilon_{y0} = 0.92 \) after 1 synchrotron oscillation. Again we find a result which seems to contradict Fig. 9. For a simulation including lattice nonlinearities and space charge but keeping the longitudinal motion frozen we again find the same results as in the previous case. This allows us to conclude that the synchrotron motion is responsible for particle extraction. The difference of beam loss between the cases with and without frozen longitudinal motion suggests that 4% of lost particles are attributable to a periodic crossing of a resonance. The resonance which appears to excite the trapping extraction process could be the 4th order \( 2Q_x + 2Q_y = 63 \). This resonance, as shown in Fig. 8c, is located right above the reference working point. The simulations performed in Fig. 9 show also that a small saturating beam loss occurs after the first synchrotron oscillation. For the case A) and B) a 0.5% beam loss is detected during \( 10^5 \) turns showing no dependence of long term beam loss on bunch maximum tuneshift.

CONCLUSION

The experimental study in the PS and comparison with simulations has reached a reasonable agreement in the emittance growth regime. With respect to previous studies we now also succeeded to confirm qualitatively the beam loss regime, even if the loss prediction is a factor 3 smaller. Further steps will be needed to get better quantitative agreement with observed loss, like adding chromaticity, more refined measurements of the nonlinearities of the used working points, etc.. As it is desirable to confirm that the modeling is accelerator independent we have started to prepare a similar measurement campaign in the SIS18 at GSI. We conclude also that for the SIS100 reference working point the beam loss budget is expected to be the extrapolated 3.6% reported in Table 3 plus \( \sim 4\% \) induced by space charge effects. More accurate calculations require the complete set of multipolar components \( b_n, a_n \) along the ring and misalignment tolerances as well as an increased set of test particles. The reliability of these beam loss predictions has to be tested in dedicated benchmarking experiments.

REFERENCES

[1] W. Henning, these Proceeding.
RF AND FEEDBACK FOR BUNCH SHORTENING

J. Tuckmantel, CERN, Geneva, Switzerland

Abstract
The motivation for, possibilities and limitations of bunch shortening by an additional superconducting RF system at 1.2 GHz supplying 43 MV in LHC are examined. Possible technical realizations are examined and critical issues shown.

MOTIVATION
The luminosity of LHC can be increased by a further transversal squeezing of the beam with stronger low-beta focusing. However, for non-zero crossing angle the leading and trailing ends of the colliding bunches may miss each other. This undesired effect can be circumvented in shortening bunches correspondingly. In this context a shortening of about a factor 2 with respect to the initial option of a 2.5 eVs bunch at 16 MV with 400 MHz is desired, i.e. a bunch length in the range of slightly above 0.5 ns.

HOW TO GET SHORTER BUNCHES
Tests in the SPS have shown that bunches as small as 0.6-0.7 eVs can be brought up to the SPS flat-top [1] from where they will be injected into LHC. There is confidence that these bunches can be captured with an emittance not larger than 1 eVs in LHC; later, after a better understanding and study of the beam transfer process, 0.7 eVs might become possible.

For an acceptable luminosity life-time due to intra beam scattering, the emittance in coast should be about 2.5 eVs [2]; a more recent study [3] suggests that an emittance of 1.75 eVs might be an acceptable option. The desired emittance will be reached by smooth blow-up of bunches using controlled RF noise during the energy ramp of LHC.

WHICH RF SYSTEM TO CHOOSE

Fig. 1a: A hypothetical bunch of 0.7 eVs at 7 TeV/c with 16 MV at 400 MHz, bunch length 0.55 ns. (Red line: bucket limit, green line RF voltage, range ±45 MV)

Fig. 2a: Standard LHC coast: a 2.5 eVs bunch with the standard 400 MHz system at 16 MV (maximum), bunch length 1.08 ns. (green line RF voltage, range ±45 MV)

Bunches of 0.7 eVs (1 eVs) will have at 7 TeV/c with the standard RF setting, i.e. only the 400 MHz system at the maximum voltage of 16 MV, just the desired bunch-length of 0.55 ns (0.66 ns). Therefore an alternative option to be studied are means to overcome the intra beam scattering limit and possible instabilities for bunches of 0.7 eVs during and towards coast. If this option is viable, the investment for the additional RF for bunch shortening is not necessary and can be spent for hardware realizing the other option.

In the following we will not consider this path but study means to shorten blown-up bunches by an additional RF system.
An increase of the total slope of the RF voltage will shorten bunches. A straightforward way consists in increasing the voltage of the existing RF system. However, there is the 4th power ‘law’ (strictly valid only for short bunches): to decrease the bunch length by a factor 2 will ask for 16 times the RF voltage. Therefore this path is unrealistic, not only for financial considerations but also due to the extreme tightness of space in the LHC tunnel and the underground areas.

Therefore a higher frequency, giving correspondingly more slope for the same voltage increase, is mandatory. The upper limit of the frequency choice is constrained by:

- the RF wavelength has to be somewhat longer than the expected bunchlength
- the cut-off tube diameter has to be at least the standard tube diameter of 5 cm, else the cavity might scrape the beam; in the RF section the standard tube diameter will be 10 cm, hence even this limit should be respected.
- to allow RF manipulations (as RF transfer) together with the ‘main’ RF system at 400 MHz, the frequency should be an integer multiple of the latter.

These constraints lead naturally to a single choice of 1.2 GHz.

To obtain the desired bunch length for a 1.75 eVs bunch, 43 MV are necessary at 1.2 GHz, working in parallel the 400 MHz system at its maximum of 16 MV. As can be seen in Fig. 4a, the nominal emittance fits into the bucket but, compared to the standard case in Fig. 2a, the margin for tails is much tighter.

If we would try to hold a standard bunch of 2.5 eVs in this configuration, the bunch would just fit the bucket with nearly no more margin at all (Fig. 4b).

Ramping up to 7 TeV/c

As to be seen in Fig. 3a, the length of a 1.75 eVs bunch ramped up with the standard 400 MHz system at 16 MV is 0.89 ns long. An oscillation at 1.2 GHz is 0.833 ns long, hence is shorter than this bunch length. Therefore one cannot ramp up bunches to the flat-top as foreseen for
standard LHC and only then transfer to the 1.2 GHz system, but the transfer has to be done earlier. On the other hand we do not want to create very short bunches at lower machine energy (the intra beam scattering life time is no true issue yet), hence the transfer should be done before or during the ramp without passing instability limits.

7 TeV, 2.5eVs, 16 MV @ 400 MHz + 43 MV @ 1.2 GHz

\[ \varepsilon = 2.503 \text{ eVs} \pm 1356.12 \text{ MeV} \ 0.72 \text{ ns} \]

Fig 4b: A 2.5 eVs bunch of 0.72 ns at 7 TeV/c with a 400 MHz system at 16 MV and 1.2 GHz system at 43 MV (green line RF voltage, range ±45 MV; the true bunch height is ±2440 MeV, the number in the plot corresponds to the ‘side lobes’ with identical Hamiltonian)

LHC with 400 MHz and 1.2 GHz system

\[ \varepsilon = 1.002 \text{ eVs} \pm 387.96 \text{ MeV} \ 1.76 \text{ ns} \]

Fig. 5a: Bunches of 1 eVs as injected and captured in LHC with 8 MV at 400 MHz, 0 MV at 1.2 GHz (green line RF voltage, range ±45 MV).

This is in fact possible using the following procedure. We can assume that after injection of all batches (and possible filamentation) we have bunches of not more than 1 eVs stored at flat-bottom with the 400 MHz system at 8 MV. This capture voltage leaves enough reserve for the RF system to damp injection oscillations and also better matching bunches coming from the SPS 200 MHz RF system. The 1.2 GHz system is forced to zero volts at this stage.

Then – before acceleration – in parallel we rise the 400 MHz system to its maximum of 16 MV and the 1.2 GHz system to about 5 MV, using it here as higher harmonic system which guarantees increased stability during the ramp for the now slightly shorter bunches compared to the standard ramp.

LHC with 400 MHz and 1.2 GHz system

\[ \varepsilon = 1.038 \text{ eVs} \pm 500.48 \text{ MeV} \ 1.50 \text{ ns} \]

Fig. 5b: 16 MV at 400 MHz, 5 MV at 1.2 GHz just before start of acceleration.

Then we accelerate, causing bunches to get shorter, to about 4.5 TeV/s where the bunch length is short enough for a gradual transfer to the 1.2 GHz system. While acceleration continues up to 7 TeV/c, in parallel the 1.2 GHz system is ramped to 43 MV and a smooth blow-up by controlled RF noise increases the bunch length so that no stability limit is crossed. The speed of blow-up can be chosen such that the desired emittance is attained simultaneously with the 7 TeV/c energy or – as in our example – blow-up may continue for a short while on flat-top till the desired emittance is obtained.

LHC with 400 MHz and 1.2 GHz system

\[ \varepsilon = 1.016 \text{ eVs} \pm 995.13 \text{ MeV} \ 0.68 \text{ ns} \]

Fig. 5c: At 4.5 TeV/c with 16 MV at 400 MHz, 5 MV at 1.2 GHz, ready for transfer to 1.2 GHz and start of blow-up.
This system alone cannot hold a (blown-up) 4 eVs bunch at injection, but one can ramp the beam up to 7 TeV/c as foreseen for the standard case, and then in parallel ramp down the 400 MHz system and continue the blow-up from 2.5 to 4 eVs, obtaining a situation as shown in Fig. 6. The ratio of bunch length to bucket length is about the same as for the standard case so that no instabilities have to be expected and on top of it the 400 MHz system works as higher harmonic system.

**BUNCH LENGTHENING OPTION**

(PIWINSKI)

Under different conditions and increased crossing angle [3], bunch lengthening can be advantageous. We will have a short look what can be done to obtain this.

A 200 MHz system with 3 MV total voltage (4 copper cavities) was designed to capture the beam from the SPS without losses [4]. Meanwhile the emittance obtained on the flat-top of the SPS for nominal intensity could be kept significantly below 1 eVs and simulations showed that attenuation of capture oscillations can be done – at the limit – with the 400 MHz system alone. Therefore, under budget constraints, it was accepted to stage the 200 MHz system.

However, exactly such a system with its 3 MV would be sufficient to run LHC in coast with longer bunches of 4 eVs emittance, hence largely above the levels to be worried concerning intra beam scattering and still leaving much (phase-) space for possible tails (see Fig. 6),

**PARAMETERS OF THE 1.2 GHz SYSTEM**

The new system should also respect the rule that there are two independent RF systems, one for each beam, to exclude beam-to-beam coupling over induced cavity fields.

As for the 400 MHz system, cavities should be superconducting. The main argument is less energy saving – as was the case for LEP with its huge total RF voltage – but the ability to create the desired voltage with a minimum of cavities having large openings (hence lower R/Q for higher order modes), resulting in a lower additional machine impedance due to the added RF system. Furthermore the presence of the beam gaps in the LHC beam forces half-detuning (i.e. the reactive beam loading compensation is only respected by a factor 1/2) so that – neglecting transients – the RF power with or without a bunch-train passing the cavity is the same, hence minimizing the peak RF power, for klystrons the decisive power.

The cavity detuning for the reactive beam loading compensation scales directly with R/Q and inversely with the operating cavity voltage, hence superconducting cavities with relatively low R/Q but high voltage need much less detuning. This detuning lets the complex RF vector drift away in one or the other direction and the RF transmitter has to send power keeping it stable as desired. Therefore a superconducting system needs less RF power to keep the beam under control.
The 400 MHz superconducting LHC cavities [5] were designed for high beam current, ‘tuning’ higher order modes to lower R/Q when possible and good mode coupling into the beam-tubes where the power and HOM couplers are placed. Also a coupling scheme with couplers exists.

Therefore it is not necessary to repeat the same optimization process for identical goals – including the particularities of superconducting cavities – but we can take over the cavity shape and its HOM couplers scaling by 1:3. A few non-scalable details will be treated later.

If we assume a maximum field level of 25 MV/m (see later), 14 cavities at 3.1 MV could do the job at the limit, so let us assume we use 16 cavities – a number also easier to subdivide in identical hardware sub-units – at 2.7 MV.

\( Q_{\text{ext}} \) for the power coupler

As for the 400 MHz cavities the requirements during the LHC cycle are too different as to work with a once and for all fixed power coupler, hence a variable power coupler is necessary. At injection the 1.2 GHz system should be as ‘invisible’ as possible hence the lowest possible \( Q_{\text{ext}} \) (even with vector feedback) is desirable.

Of course one might think at injection to detune the cavity far away from the operating frequency and even run without power then. This reduces the induced RF power certainly – due to beam gaps and bunch inequalities there is some RF current ‘everywhere’ – but the cavity impedance is always somewhere close to another revolution frequency multiple, these being very dense due to the low revolution frequency of about 11.2 kHz. In any case, at one moment the cavities have to approach the nominal tune to go into operation.

As for the 400 MHz cavities, the minimum \( Q_{\text{ext}} \) obtainable is about 10,000, else the antenna penetrates the volume where the beam passes. The maximum stroke of the variable coupler is 20 there, hence a maximum of 200,000 should be obtainable. As we shall see later, this range covers the desired optimum values.

HOM impedances

The R/Q of longitudinal modes remains identical when scaling the cavity dimensions, hence the fundamental mode at \( f=1.2 \text{ GHz} \) has \( R/Q=45 \Omega \); the same identity scaling holds for all longitudinal HOMs.

We scale the main HOMs calculated for the 400 MHz cavity [6] and assume damping as in [7]. Then a total impedance (excluding the fundamental mode) of less than 10 k\( \Omega \) can be accounted for recognizable longitudinal modes up to 3.5 GHz. Evidently there will be a lot more ‘little’ modes in higher frequencies, as usual impossible to identify clearly, partly propagating along the beam tube.

For transversal modes dipoles are dominant, and there the field rises linearly from the axis where it is zero, hence a particle with the identical off-axis trajectory induces a 3 times higher longitudinal voltage and hence a probe particle on the same off-axis trajectory feels a 3 times higher deflection in the scaled cavity.

Fundamental mode impedance

Assuming the worst case, the maximum \( Q_{\text{ext}} \) of 200,000, the bare impedance per cavity is 9 M\( \Omega \) or 144 M\( \Omega \) for the whole system of 16 cavities (always per beam). It appears evident that (as for the 400 MHz cavities) the 1.2 GHz cavities need an individual fast vector feedback to reduce this impedance and that therefore single cell cavities have to be used to avoid pass-band problems in the feedback loops. The loop delay, mainly due to the inevitable layout in the tunnel for radiation protection of sensitive equipment, is at least \( T=600 \text{ ns} \). The loop gain where the system starts to oscillate on its own (see e.g. [8]) is \( g_{\text{osc}} = Q_{\text{ext}}/(2f_{RF}T) \) and we have to stay away from it by at least a factor 2. Therefore the on-tune impedance with feedback becomes in the best of all cases \( Z_{\text{RF}} = 4(R/Q)f_{RF}T \) or 130 k\( \Omega \) per cavity, hence 2 M\( \Omega \) for the whole system (per beam).

Necessary RF Power

As mentioned in [3] we consider 3 scenarios for the DC beam current: (n)ominal 0.56 A, (u)ltimate 0.85 A and (s)uper 1.41 A. The bunch length is about 15 cm in all cases and (assuming a cos\(^2\) longitudinal profile) the RF current at 1.2 GHz for such bunches is reduced to about a factor of 0.78 compared to point bunches. To minimize the necessary RF power for a beam with gaps – but neglecting RF transients – we have to half-detune the cavities. Under these conditions we obtain the following minimum RF power values at 1.2 GHz of Tab. 1. It should not be forgotten that in the real machine there are non-negligible RF transient power spikes, tuning and \( Q_{\text{ext}} \) deviations from the optimum values and small phase deviations between cavities so that some cavities give power to the beam which is taken back (and dumped) by others.

Tab. 1: Minimum power to keep zero volt on a tuned cavity for \( Q_{\text{ext,min}} \) and to create 2.7 MV assuming the optimum coupling \( Q_{\text{ext,opt}} \) and perfect half-detuning. Any (non negligible) RF transients and setting errors are not included and extra power has to be provided for it.

<table>
<thead>
<tr>
<th>DC current</th>
<th>( P(0) ) [kW]</th>
<th>( Q_{\text{ext,min}} ) on tune</th>
<th>( Q_{\text{ext,opt}} )</th>
<th>( P(2.7) ) [kW]</th>
</tr>
</thead>
<tbody>
<tr>
<td>n (0.56 A)</td>
<td>43</td>
<td>135000</td>
<td>300</td>
<td></td>
</tr>
<tr>
<td>u (0.85 A)</td>
<td>100</td>
<td>90000</td>
<td>450</td>
<td></td>
</tr>
<tr>
<td>s (1.41 A)</td>
<td>275</td>
<td>55000</td>
<td>750</td>
<td></td>
</tr>
</tbody>
</table>

There exist one further possible option in the far future. The principal method is under theoretical and simulation study for the 400 MHz RF system [9], but the method cannot be relied on before it has been validated by a real test with LHC once the machine runs and MDs of such type will be accepted by the physics experiments.

If there would be no beam instabilities, we might simply supply each cavity with sufficient but constant RF power keeping the average field up (but not compensating beam loading), the beam taking no energy...
in coast, and adjust the detuning for minimum power. Bunches, in coast riding at the RF zero crossing, would then 'slide to' an individualized position slightly different from the regular bunch spacing (the deviation is very small and would not harm the physics experiments [10]) and – due to uncompensated beam loading and cavity detuning – the phase of the voltage vector would slightly swing forward (bunch train) or backward (beam gap). Evidently beam loading has to be low enough, the cavity inertia (filling time and stored energy) high enough and the beam-gaps sufficiently distributed so that the voltage does not vary very much. This is the case here and the beam would coast with a minimum of RF power and no transient power spikes.

However, under these conditions in reality multi bunch instabilities would develop rapidly, the RF voltage has to be kept under tight control. This can in theory be done if the 'set-values' of the fast RF feedback along a machine turn could be set to exactly prescribe what the voltage did in the case depicted above, i.e. we would have the same voltage and bunch positions as above but now under tight feedback control. Then any developing instability will show up as deviation from these set-values and the RF counteracts immediately guaranteeing beam stability; the impedance is reduced as above but now under tight feedback control. Then any developing instability will show up as deviation from these set-values and the RF counteracts immediately guaranteeing beam stability; the impedance is reduced as in the standard case.

The practical obstacle is to establish the correct 'set-values' which depend in a complex way on the never perfect cavity settings and scattered beam parameters, especially the bunch charges, all not precisely known nor exactly measurable. Several methods to determine 'set-values' without direct knowledge of the beam and cavity parameters are under study. All methods do not exploit the knowledge of (in the simulation program evidently present but in reality unavailable data. All lead – more or less rapid – to the desired result of nearly constant and relatively low RF power demand. An example is shown in [10], where the initial spikes, while forcing bunches onto their nominal position (Fig. 13a) disappear (Fig. 13b); the 'set-values' and the nearly flat power curve are also shown in Fig. 13b.

The open question remains if these methods and the corresponding hardware are robust enough to cope with real life imperfections, considering that we 'play with' a high gain feedback loop and each error causes a reaction multiplied by the gain; hence we have to wait for a real test with LHC before relying on such a, on paper, very attractive method.

If the method could be applied, for a maximum setting of $Q_{\text{max}}=200,000$ and after establishing the 'set-values', a power of about 100 kW might be sufficient for all beam current cases. In the intermediate state to establishing the 'set-values', a higher power level will be necessary, the size of the latter being difficult to estimate before real tests are done.

### SC. cavity technology

The LHC 400 MHz cavities are operating at 5.5 MV/m, a field level where the CERN developed Nb/Cu sputter technique gives stable and still high Q cavities.

For the 1.2 GHz cavities we have assumed fields of 25 MV/m. In fact this is not an empty hypothesis, but in the framework of the TESLA project [11] this performance has been established regularly in a real machine (FEL). Even considerably higher values were obtained nearly regularly but since a 'lost' cavity in LHC causes a complete loss of the beam (the beam-dump has to be triggered to avoid component damage [12]), we prefer to stay conservative in this framework.

TESLA works at 1.3 GHz which is very close our frequency of 1.2 GHz. Cavities are made from bulk-niobium of high quality, achieving $Q_0$-values larger $10^{10}$. The operating temperature is 2 K, not 4.5 K as for the 400 MHz cavities, but this technology is used for cavities also at several other laboratories (e.g. CEBAF/TJLAB) and also the LHC magnets work – under slightly different conditions – at 1.8 K.

One difference is that TESLA is a pulsed machine with 1.4µs RF-on time per pulse but a repetition rate of only 5 (or 10) Hz, i.e. a very low duty cycle, while LHC asks for CW operation. The total liquid helium consumption in CW operation for the many km long TESLA would be excessive, but the section in LHC is very modest so that the total consumption (for 2x16 cavities at 2.7 MV about 250 W dynamic losses at 2 K) is still manageable but not negligible at all, adding also static losses (see later). Another problem might be the cryogenic liquid throughput since the whole 1.2 GHz system might be advantageously housed in a single or two cryostats.

### Scaling Considerations

We have assumed that we can scale the 400 MHz cavity by 1:3 to get an operational 1.2 GHz cavity. This holds for most details but not all.

#### The Cavity

Fig. 6 shows a 400 MHz cavity and in the same scale three 1.2 GHz cavities together with some dimensions. The cut-off tube of the 1.2 GHz cavity scales to 10 cm which is acceptable, identical to the standard tube diameter in the RF section. But this means that there cannot be any taper as for the 400 MHz cavities (with a diameter reduction of 3:1) and higher frequency HOMs may propagate along the whole RF section.

A more serious problem are the Conflat® flanges on all cavity ports that have to guarantee ultra high vacuum tightness. The 400 MHz cavities use M6 bolts which would scale to M2 bolts and there are some doubts that they are sufficient for this purpose. Therefore relatively larger flanges with larger bolts had to be used which is not trivial in view of the many close ports.

Another non-scalable quantity is the cavity wall thickness, having to resist the liquid helium pressure...
against vacuum and being adequate for the manufacturing process. For the same relative Δf-tune stroke, obtained by elastic longitudinal cavity deformation, the same relative movement is necessary, hence elastic limits – already close for the 400 MHz cavities – have to be checked as well as the necessary strength of the (also scaled) tuning frame and motor drive.

![Image of 400 MHz cavity in Nb/Cu technology](image)

Fig. 7: Virtual reality of a 400 MHz cavity in Nb/Cu technology (top) and three 1.2 GHz cavities (bottom) in the same scale with some dimensions. The 400 MHz cavities do exist in reality.

**HOM coupler and ports**

The HOM port diameter is also reduced by a factor 3 which would result in an at least (skin effect) 9 times higher power density. However, the wave length of these modes is generally larger than the bunch length so that the absolute power coupled out will we reduced correspondingly. Therefore this point should be of no serious concern.

The (superconducting) HOM couplers of hook-type are shorter by a factor 3, hence conduction cooling is improved. The hook is filled with liquid helium and a small tube is introduced to let evaporated helium gas escape. A study has to clarify if a hollow hook with inlet tube can still be produced at reasonable cost (watchmaker’s work) or if conduction cooling is already sufficient.

**Power coupler and port**

The main problem for this RF system is the power coupler. From Tab. 1 we see that even for the lowest beam current (nominal) 300 kW are required as absolute minimum power; this means that in the order of 400 kW (or more) should be foreseen to handle transients and setting imperfections. The 400 MHz couplers are designed to handle 300 kW nominal – which they do now on the test stand after serious development work – but they have to be fabricated, treated and baked out with utmost care, and then RF processed for a long time.

The 400 MHz couplers have an outer conductor diameter of 14.5 cm, the scaled 1.2 GHz version then about 5 cm. This means that for the same power flow the field strength is 3 times higher; one should always keep in mind that sc. cavities with feedback may completely reflect the incident wave, creating local fields twice as high. The local power density, taking also the change of skin-depth with frequency into account, is increased by a factor of about 15 compared to the 400 MHz coupler. For the higher beam current options things are even worse (see power data of Tab. 1).

**Power transmitters**

A further bottleneck is the power transmitter. In theory one can always add the power of several smaller transmitters but this method becomes prohibitive for more than doubling and system complexity and cost increases correspondingly.

As example, compiled by Erk Jensen, CERN-AB, two high power transmitters:

- The TESLA klystrons at 1.3 GHz (designed for pulsed operation) can deliver ‘only’ 150 kW in CW.
- Raytheon has developed the "Amplitron" that delivers 425 kW CW but at 3 GHz, which is more than twice the target frequency and scaling to 1.2 GHz is difficult.

It should be reminded that for the highest beam current (s) also the 400 MHz system with 300 kW – already for the (u) option stretched to its limit – would need an upgrade of klystrons, circulators and loads. If the existing power coupler can still be used has to be found out.

Therefore also there an R&D effort in collaboration with or only by industry is required.

**The cryostat(s)**

The 400 MHz cavities are so large that at a given position one cavity can be placed, accelerating one beam and letting the counter-rotating beam pass close to it (in a tube still housed in the cryostat, see Fig. 8). Therefore the 400 MHz modules are placed one after the other along the ring in staggered position so that the two centre modules accelerate one beam and let the counter-rotating one pass next to it, the outer two modules doing the complement.

The 1.2 GHz cavities are small enough that two of them can be placed at the same longitudinal position with an axial distance of 42 cm so that both beams can be accelerated there, shortening the necessary length in the ring, a precious commodity. This allows housing all cavities in a single cryostat (Fig. 9a) or in two parallel cryostats (Fig. 9b).

Both cryostats are not easy to be designed. First they are small and cryogenic insulation difficult to maintain. But most of all there is a huge quantity of inlets and outlets that all present heat leaks and – even more trivial – space has to be allocated so that also wave-guides or cables can be mechanically connected. To remind, 16 such cavities will need 16 power couplers of 300 kW (or more) capability, 16 tuner with mechanical gear below (?) the cryostat, 4x16=64 HOM outlets with connectors, a liquid helium inlet, a gaseous helium outlet, vacuum connections for beam and insulation vacuum, vacuum and level gauges. The double cryostat of Fig. 9a will need all RF connections twice, the others only once. Such a cryostat will certainly ask the effort of a good cryogenics engineer.
CONCLUSION

LHC bunches can be shortened as desired and brought up to 7 TeV/c by adding a 1.2 GHz RF system at 43 MV to the existing 400 MHz system at 16 MV. To realize such a system, three groups of items can be classified:

1) Ready to be used are the cavity shape and HOM damping system (scale 400 MHz system), the cavity technology (copy TESLA technology) and cryogenic supply (TESLA, CEBAF and LHC). The necessary tunnel length of 15-20 m should be available.

2) Engineering efforts have to be invested into the detailed cryostat design with its many RF in- and outlets and the allocation of the waveguides and power transmitters in the tight underground space.

3) Possible show-stoppers, asking for a considerable R&D effort, are the availability of very high power RF transmitters at 1.2 GHz and, most of all, a power coupler compatible with superconducting cavities that can carry much more than 300 kW on an outer coaxial tube diameter of about 5 cm without overheating, arcing or multipacting.

Furthermore, beam stability issues concerning the ramping and coasting scenario should be verified as well the possibility to use 1.75 eVs bunches, instead of 2.5 eVs, concerning the intra beam scattering limit.

Finally, some thought should be spent to examine the possibility of directly fighting intra beam scattering and its undesired blow-up effect. In case of a positive reply the necessary hardware could be paid by the money saved from the then not necessary 1.2 GHz RF system, provided the technical difficulties are less.
References

and V. Rödel and L. Verolino, Geometry of a superconducting 400 MHz accelerating cell for LHC, CERN-SL/RFS/Note 91-11
[9] J. Tuckmantel, Adapting feedback set-values to reduce power spikes in the LHC RF system, to be published.
ELECTRON CLOUD EFFECTS – OBSERVATIONS, MITIGATION MEASURES, AND CHALLENGES IN RHIC AND SNS

J. Wei, M. Blaskiewicz, W. Fischer, H.C. Hseuh, U. Iriso, T. Roser, L. Wang, S.Y. Zhang
Brookhaven National Laboratory, New York, USA

Abstract

Electron cloud is one of the leading mechanisms that limit the performance of high intensity circular accelerators and colliders. In the Relativistic Heavy Ion Collider, multi-bunch electron cloud effects are observed both in the warm region and superconducting region when the number of ion bunches and their intensities are raised beyond the design values. Vacuum-pressure rises, transverse tune shifts, and electron detector signals are observed at injection, upon transition crossing, and at top energy. Transverse emittance growth, fast instabilities, and beam loss also occur upon transition crossing. With the Spallation Neutron Source Ring, single-bunch electron cloud effects are expected for the high intensity proton beam. A comprehensive list of mitigation measures are implemented both to reduce the production of electron cloud and to control the beam stability. This paper intends to provide an overview of observations, performance limitations, and beam dynamics challenges pertaining to electron cloud build-up in high intensity, circular hadron accelerators.

1 INTRODUCTION

Since the first reports four decades ago, electron cloud has attracted the attention of accelerator physicists around the world [1]-[5]. During the recent years, electron-cloud effects are found to limit the performance of several high-intensity and high-brightness circular accelerators [6]. A fast, transverse electron-proton instability and the induced beam loss limits the beam intensity in the Proton Storage Ring (PSR) at the Los Alamos National Laboratory (LANL) [7]. Transverse emittance blow-ups caused by the electron cloud limit the luminosity in the lepton factories (BEPC, KEKB, PEP-II) [8]-[11]. During recent operations of the Relativistic Heavy Ion Collider (RHIC), vacuum-pressure rises associated with electron-induced gas desorption are found to limit the number of stored bunches and the beam intensity during high-intensity operations (Fig. 1) [12]-[14]. Electron-cloud phenomena include transverse tune shifts (KEKB, AGS Booster, RHIC), coupled-bunch (B factories, BEPC, PS, SPS) and single-bunch (KEKB, SPS, PSR, RHIC) instabilities, vacuum-pressure rise (RHIC etc.), emittance growth (KEKB, PEP-II, SPS, RHIC), beam diagnostics interference (RHIC, PS, SPS, PSR), and heat load on superconducting cryogenic wall (SPS beam experiments) [15].

Electron cloud is expected to be one of the performance-limiting mechanisms for the Spallation Neutron Source (SNS) accumulator ring currently under construction at the Oak Ridge National Laboratory [16]. When the full beam intensity of 2 \times 10^{14} protons per pulse is accumulated in the ring, trailing-edge electron multipacting is expected to occur. Several mitigation strategies are adopted in the design, both in suppressing electron-cloud formation and in enhancing Landau damping.

This paper mainly discusses electron-cloud effects in RHIC and SNS. Section 2 briefly reviews key mechanisms of electron-cloud formation. Section 3 summarizes the experimental observations in the RHIC collider. Section 4 discusses theoretical expectations in the SNS accumulator ring. Mitigation measures adopted in RHIC and SNS are presented in Section 5. Beam-dynamics challenges are given in Section 6.

* Work performed under the auspices of the U.S. Department of Energy.
† jwei@bnl.gov
2 MECHANISMS

Key mechanisms pertaining to the formation of electronic cloud are beam-driven electron multipacting and electron trapping. Depending on the beam parameters, the multipacting can be classified into several regimes: multiple short-bunch multipacting [4, 17, 18, 19], single long-bunch, trailing-edge multipacting [7, 20], and the intermediate regime. For the long-bunch and dc-beam [2, 3] cases, beam-induced trapping is critical to sustain electron concentration. For the short-bunch cases, trapping due to magnetic field is suspected to be responsible for the long electron lifetime over beam gaps [21].

2.1 Multiple, short-bunch regime

The phenomena of beam-driven multiple short-bunch multipacting have been observed in many lepton accelerators. The multibunch multipacting occurs if the transit time of the electrons crossing the vacuum pipe is comparable to the time between successive bunches, and if the electrons gain enough energy to produce more than one secondary electron when they hit the vacuum-pipe wall [4].

The multiple, short-bunch regime is defined as

\[ \zeta_m \equiv \frac{2b \beta}{s_b \beta_e} \geq 1 \]  

(1)

where the short-bunch multipacting parameter \( \zeta_m \) is defined as the ratio between the transit time of the characteristic electron crossing the vacuum pipe to the time between successive bunches, \( s_b \) is the distance between the subsequent bunches, \( b \) is the average radius of vacuum pipe, \( \beta_e \) is the velocity of the beam, and \( \beta_e \) is the characteristic velocity of the electrons [4, 22]. In the case that \( \zeta_m \approx 1 \), the energy gained by the electron from the passage of the beam bunch is approximately given by

\[ \Delta E_e = m_e c^2 \left( \sqrt{\left( \frac{2r_e N_0}{\beta b} \right)^2 + 1} - 1 \right) \approx 2m_e c^2 \left( \frac{r_e N_0}{\beta b} \right)^2 \]  

(2)

where \( r_e = e^2 / 4\pi \varepsilon_0 m_e c^2 \) is the classical radius of electron, \( N_0 \) is the number of particle in the beam bunch, the second relation is true if the electron motion is non-relativistic, and

\[ \zeta_m \approx \frac{\beta \beta_e^2}{r_e N_0 s_b} \]  

(3)

In the general case that \( \zeta_m > 1 \), the energy gain needs to be analyzed numerically [23]. In order to sustain a multipacting, the energy gained by the electron upon one or more beam bunch passage must be such that the overall electron-induced secondary-emission yield \( Y_{ee} \) (SEY) satisfies

\[ Y_{ee} > 1 \]  

(4)

Multibunch electron multipacting may occur for almost any value of \( \zeta_m \) satisfying Eq. 4 [8].

The occurrence of electron-cloud multipacting build-up depends strongly on the beam intensity and the secondary-emission yield of the beam-pipe surfaces. The multiple short-bunch multipacting is sensitive to the bunch spacing and bunch pattern, and insensitive to the bunch length and profiles when the condition of Eq. 1 is satisfied. The effect is usually stronger for regions of smaller vacuum-pipe aperture where the beam field at the wall and the electron crossing frequency are both high. Both multi-bunch and single-bunch head-tail instabilities are observed. Mechanisms that sustain electron concentration include trapping by the magnetic fields (e.g., quadrupole-field trapping resonant to the bunch passage [21]) and near-elastic reflection of low-energy electrons [24].

2.2 Single, long-bunch regime

The phenomena of beam-driven single, long-bunch multipacting are observed in routine operations in the PSR and expected in the SNS ring. Single-bunch, trailing-edge multipacting starts to dominate if the bunch length is long enough to sustain multiple passes of electrons. As shown in Fig. 2, electrons are attracted by the rising beam particle density as the beam potential increases [25, 26]. The motion is characterized by the electron bounce-frequency

\[ \omega_e = c \sqrt{2 \pi r_e n_p} \]  

(5)

where \( n_p \) is the volume density of the beam. After the passage of the beam density peak, electrons are released and accelerated by the part of the beam of decreasing density (trailing edge of the beam density distribution). The number of electrons grows dramatically upon such trailing-edge multipacting, as observed at the PSR [7].

The single, long-bunch regime is defined as

\[ \zeta_s \equiv \frac{4b \beta}{s_b B_f \beta_e} \ll 1 \]  

(6)

where the long-bunch multipacting parameter \( \zeta_s \) is defined as the ratio between the transit time of the characteristic electron crossing the vacuum pipe to the passage time of
When the electron motion is non-relativistic, the energy gained by an electron is approximately given by

$$\Delta E_e \approx 4m_e c^2 \beta_b \frac{\tau_e N_0}{s_I B_f}$$

while

$$\zeta_e \approx \frac{\beta_b}{\sqrt{\tau_e N_0 s_I B_f}}.$$  \hspace{1cm} (8)$$

Single-bunch multipacting occurs if the energy gained by the electron is such that

$$Y_{ee} > 1.$$  \hspace{1cm} (9)$$

Fig. 3 shows examples of the measured secondary yields as functions of the primary electron energy on vacuum-chamber surfaces of the SNS ring.

Electron-cloud build-up due to the single-bunch multipacting is typically not sensitive to the bunch spacing, while electron survival between the bunch passage is sensitive to the trapping in the beam gap. The build-up depends critically on the length of the beam bunch and the variations in its longitudinal density, which determine the energy gain and the multipacting duration. The effect can be stronger for regions of larger vacuum-pipe aperture where higher energy gained over a longer path-length results in a higher secondary yield, although the longer path-length may also reduce the multipacting frequency. Upon acceleration by the beam, the electron energy is typically below kV level. Associated with the electron-cloud build-up, single-bunch, fast transverse instabilities are observed [7]. Main trapping mechanisms include attraction of the electrons due to

the increasing beam density at the rising density edge of the bunch, reflections from the vacuum-chamber wall, and trapping by the beam residual in the gap. Trapping by the magnetic fields (e.g., quadrupole field) is typically weak due to lack of resonance mechanism [25]).

A coasting, dc beam is a special case of the long-bunch regime. In this case, the electrons are trapped by the constant beam potential, and the electron density may accumulatively increase leading to vacuum pressure rise and two-stream instabilities when the thresholds are exceeded [2, 3].

### 2.3 Intermediate regime

RHIC belongs to the intermediate regime where the transit time of the electrons crossing the vacuum pipe is comparable to the bunch length, as shown in Fig. 4. Between the subsequent bunches, the electrons typically reflect from the vacuum chamber wall for several times.

Upon acceleration by the beam bunch, the electrons gain energy up to several hundred eV, and hit the wall with an average SEY $Y_{ee,0}$ typically much larger than 1. Upon subsequent impacts on the wall, the electron energy is low (typically below 10 eV) due to lack of beam potential, and the average secondary-emission yield $Y_{ee,i}$ ($i = 1, \ldots$) is smaller than 1. We define the characteristic SEY $Y_{ee,C}$ as the product of the average SEY of $N_{ee}$ reflections between two subsequent bunch passage. The threshold for the multipacting in this regime corresponds to the condition when

$$Y_{ee,C} \equiv \prod_{i=0,1,\ldots,N_{ee}} Y_{ee,i} > 1$$

where

$$Y_{ee,0} > 1, \text{ and } Y_{ee,i} < 1 \text{ for } i = 1, \ldots, N_{ee}$$

Multipacting in the intermediate regime shares the features of both the short-bunch and long-bunch regimes. The build-up is not only sensitive to the bunch spacing and bunch pattern, but also sensitive to the bunch length and peak intensity. With a shorter bunch spacing, the number of low-energy electron passages $N_{ee}$ is reduced, increasing
the chance of exceeding multipacting threshold \( Y_{ee,C} > 1 \). With a shorter bunch length and higher peak beam intensity, the electron energy gain upon bunch acceleration becomes higher, leading to a higher secondary emission yield \( Y_{e,0} \) and \( Y_{ee,C} \). Mechanisms that sustain electron concentration include reflections from the vacuum-chamber wall, trapping by the beam residual in the gap, trapping by the magnetic fields, and possible secondary ionization and trapping [28].

In the following section, we present observations pertaining to electron cloud formation in RHIC.

## 3 RHIC OBSERVATIONS

The RHIC collider is comprised of two quasi-circular rings (blue ring and yellow ring) each consisting of superconducting magnets contained in 4 K cryostats (Table 1). Counter-circulating ions of various species from proton to gold can be stored in each ring to intersect with each other at up to six room-temperature locations.

### Table 1: Typical machine and beam parameters pertaining to electron cloud in RHIC.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ring circumference</td>
<td>3833.8 m</td>
</tr>
<tr>
<td>Aperture, IR (2/6/8/10, 4/12)</td>
<td>7, 12 cm</td>
</tr>
<tr>
<td>Aperture (arc, triplet)</td>
<td>7, 13 cm</td>
</tr>
<tr>
<td>Beam species</td>
<td>p - Au</td>
</tr>
<tr>
<td>Energy, injection - top, Au</td>
<td>9.8 - 100 GeV/u</td>
</tr>
<tr>
<td>Energy, injection - top, p</td>
<td>24.3 - 100 GeV</td>
</tr>
<tr>
<td>Transition energy, ( \gamma_T )</td>
<td>22.9</td>
</tr>
<tr>
<td>Rebucketing energy, ( \gamma_{top} )</td>
<td>107.5</td>
</tr>
<tr>
<td>Bunch intensity, p, Au</td>
<td>( 10^{11}, 10^9 )</td>
</tr>
<tr>
<td>Bunch spacing</td>
<td>108, 216 ns</td>
</tr>
<tr>
<td>Bunch length, full</td>
<td>35 - 5 ns</td>
</tr>
<tr>
<td>Electron spacing</td>
<td>150 - 400 MHz</td>
</tr>
<tr>
<td>Peak bunch potential</td>
<td>0.25 - 1.6 kV</td>
</tr>
<tr>
<td>( e^- ) energy gain upon acceleration</td>
<td>50 - 300 V</td>
</tr>
</tbody>
</table>

During normal operations, ion bunches at 216 ns spacing, each containing up to \( 10^{11} \) charges (e.g., \( 10^9 \) Au\(^{70+}\) or \( 10^{11} \) protons), are injected into each ring and accelerated to the top energy for hours of storage. During the acceleration cycle, the bunch length reaches the minimum, and the peak intensity reaches the maximum when the beam crosses the transition energy (\( \gamma_T \approx 22.9 \)), and when the beam is transferred from the acceleration RF system (28 MHz) to the storage RF system (200 MHz) at the beginning of the storage (rebucketing, \( \gamma_{top} \approx 107.5 \)).

Sources of seed electrons in RHIC are expected to be gas ionization and beam-loss induced surface emission and desorption. The electron cloud becomes measurable only if the threshold for intermediate-regime multipacting and trapping, as discussed in the previous section, is exceeded. Observables include vacuum pressure rise in the warm-bore region, gas density increase in the cold-bore region, electron flux on the wall, as well as bunch-dependent coherent tune shifts, transverse instabilities, transverse emittance growths, and beam losses.

### 3.1 Vacuum pressure rise and electron flux

One of the most pronounced phenomena in RHIC pertaining to electron cloud is the vacuum pressure rise when the spacing between full-intensity bunches is reduced, as shown in Fig. 1 [12]. Pressure rise often occurs during beam injection when a large number of bunches are injected with a reduced bunch spacing (108 ns in comparison to the design value of 216 ns), and upon transition crossing and rebucketing when the bunch peak intensity reaches the maximum. Figs. 5 and 6 show a direct correlation between the vacuum pressure and the integral electron flux on the wall measured by a retarding-field electron detector [14].

The pressure rise often occurs in the interaction region (IR) where beams in the two rings are both present effectively doubling the local beam intensity, in regions where unbaked surfaces are exposed to the beam, and in regions where high SEY material surfaces (e.g. Be vacuum chamber in the PHOBOS experimental region) are present. Fig. 7 shows an example of the electron flux on the IR vacuum-chamber wall measured by the electron detector, where even the intensity threshold for electron multipacting is not exceeded in each individual ring, the multipacting occurs in the interaction region. Fig. 8 shows that although a beam gap of 432 ns appears to suppress the electron multipacting, the suppression is incomplete since the multipacting re-occurs in a much shorter time [14].

In addition to the beam-induced electron multipacting, beam-induced vacuum run-away caused by gas desorption from the pipe surface and subsequent ionization by the beam may also contribute to a pressure increase. During
the year 2003 - 2004 operations, pressure rises in the unbaked regions (un-baked stochastic cooling kickers and un-baked collimator jaws) are attributed to the increase in both the electron secondary-emission yield and the gas desorption yield in comparison with that of the baked surfaces [29].

### 3.2 Tune shifts and electron-ion instability

During the RHIC acceleration cycle, instability is likely to occur upon transition crossing due to lack of synchrotron oscillation and synchrotron frequency spread [30]. Transition-energy jump is used to effectively increase the transition crossing rate. However, transverse instabilities of both slow (~100 ms) and fast (~15 ms) growth rates are still observed [31]. Some cases of transition instability are not correlated to the observation of electron cloud, and are often cured by the adjustment of the timing and magnitude of the chromaticity jump, and the activation of octupole magnets at transition [32]. Systematic studies on a fast instability at transition are in progress during year 2004 - 2005, indicating correlation to the electron cloud [33]. Fig. 9 shows the transverse coherence signal defined as the time-averaged rms amplitude of the transverse centroid displacement. A transverse instability occurs immediately after transition for about 0.1 s, leading to beam loss and emittance growth that are increasingly severe for later bunches of the bunch train, as discussed in the next section. This fast instability and beam loss also results in bunch-dependent longitudinal distributions, as shown in Fig. 10.

Bunch-dependent coherent tune shifts in both the horizontal and vertical directions have been observed during the injection of proton bunches at an intensity of $0.3 \times 10^{11}$ per bunch and 108 ns bunch spacing. The tune shift of about $2.5 \times 10^{-3}$ corresponds to an electron density of up to 2 nC/m [13]. Measurement of bunch-dependent tune shift during acceleration ramping is yet to be attempted.

### 3.3 Emittance growth and beam loss

Beam experiments during year 2004 - 2005 indicate that with the beam in a single (blue) ring at 108 ns bunch spacing of $5 \times 10^{9}$ copper ions per bunch, electron cloud ef-
Figure 9: Horizontal coherence observed on bunch #39 indicating instability after transition. 40 bunches, each of $5 \times 10^9$ copper ions, are accelerated at 108 ns spacing. The shift in baseline is due to the orbit change upon transition jump. The coherence that corresponds to an electron-ion instability causing a large beam loss, occurs within about 0.1 s immediately after transition crossing [33].

Figure 10: Longitudinal profiles of the first (top) and the 14th (bottom) bunch immediately after transition in the RHIC blue ring. The beam parameter is the same as Fig. 11 [33].

Figure 11: Measured peak electron flux on the wall as a function of time during the RHIC operation cycle. 40 bunches, each of $5 \times 10^9$ copper ions, are accelerated at 108 ns spacing. The flux increases as the beam approaches the transition energy, and disappears due to the large beam loss at transition [33].

Figure 12: Measured bunch-dependent intensity reduction near transition crossing in RHIC. The beam parameter is the same as Fig. 11. The wall-current monitor data is not logged for about 1 minute near transition [33].

Factors dominate the acceleration ramp [33]. As the beam approaches transition, the vacuum pressure and the electron flux (Fig. 11) both increase due to the increase of the peak beam intensity. On the other hand, the electron cloud impacts the beam mostly at the time of transition crossing (within about 0.1 s) when the beam particle motion is non-adiabatic, causing electron-ion instability (Fig. 9 obtained from the beam-position monitor data), beam loss (Figs. 12 and 13 obtained from the wall-current monitor data), and emittance growth (Fig. 14 obtained from the ionization pro-
Systematic studies indicate that families of octupole magnets can enhance Landau damping and effectively reduce the electron-ion instability beam loss. A lower RF voltage at transition also reduces the peak intensity and the momentum spread, effectively reducing both the chromatic [30] and the electron-cloud effects [33] at transition.

### 3.4 Background and interferences

Occasional, excessive experimental background in the interaction region is attributed to the ionization when the vacuum pressure rises due to electron cloud (Fig. 15), and to secondary shower from direct beam loss on the wall [34].

Electron cloud is also found to interfere with the beam diagnostics in RHIC. The RHIC ionization profile monitor (IPM) system is based on collecting electrons produced from controlled gas ionization with the circulating beam. Sweeping electrodes are used, and the multi-channel plates for electron collection need to be recessed from the vacuum-chamber wall to avoid superfluous signals (Fig. 16) [35].

### 3.5 Cold-region density increase

The molecular density in the 4.5 K temperature cold-bore region of the superconducting magnets has been indirectly measured with cold-cathode gauges connected through 1.5 m long, 2.5 cm diameter conduits with about 1 l/s conductance, capable of measuring the effective pressure at $10^{-10}$ Torr level [29]. As shown in Figs. 17 and 18, the molecular density increases in the cold and warm regions are similar in time scale and magnitude. An increase of up to three orders of magnitude in the effective pressure is observed both in the triplet region (Q1-Q2-Q3) and in the arc region (Q9-Q20-Q9) around the machine (Fig. 19). Temperature rise is not observed at a resolution of about 0.01 K.

Figure 13: Beam loss at transition as a function of bunch sequence number in the 40-bunch train of copper ions in the RHIC blue ring. The beam parameter is the same as Fig. 11 [33].

Figure 14: Beam horizontal emittance growth at transition as a function of bunch sequence number in the 40-bunch train in the RHIC blue ring. The beam parameter is the same as Fig. 11 [33].

Figure 15: Transition pressure rise at high beam intensity causes high experimental background in the d-Au operation in RHIC. After beam-collision starts, the pressure reduces; but the corresponding background is still 200 kHz, much higher than the coincidence rate of 7 kHz.
corresponding to a heat load of about 5 W per 100 m.

The effective pressure rise in the cold region is attributed to lack of active pumping during most machine shutdowns. An estimated 10 to 100 mono-layers of molecules are accumulated. Direct measurement of the electron flux in the cold region has not been possible, especially since all the beam-position monitors are grounded.

4 SNS EXPECTATIONS

The Spallation Neutron Source (SNS) accumulator ring is designed to accumulate, via H− injection, protons of 2 MW beam power at 1 GeV kinetic energy at a repetition rate of 60 Hz. At such beam intensity, electron-cloud is expected to be one of the intensity-limiting mechanisms that compli-
Table 2: Typical machine and beam parameters pertaining to electron cloud in the SNS accumulator ring.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ring circumference</td>
<td>248.0 m</td>
</tr>
<tr>
<td>Aperture (arc, straight)</td>
<td>21, 30 cm</td>
</tr>
<tr>
<td>Beam species</td>
<td>proton</td>
</tr>
<tr>
<td>Kinetic energy</td>
<td>1 GeV</td>
</tr>
<tr>
<td>Bunch intensity</td>
<td>$2 \times 10^{14}$</td>
</tr>
<tr>
<td>Beam gap length</td>
<td>250 ns</td>
</tr>
<tr>
<td>Bunch length, full</td>
<td>700 ns</td>
</tr>
<tr>
<td>Electron bounce frequency</td>
<td>~ 175 MHz</td>
</tr>
<tr>
<td>Peak bunch potential</td>
<td>5 - 10 kV</td>
</tr>
<tr>
<td>$e^{-}$ energy gain upon acceleration</td>
<td>0 - 300 V</td>
</tr>
</tbody>
</table>

Figure 20: Computer simulation of the beam-induced electron multipacting in the non-magnetic region of the SNS ring. The beam intensity is $2 \times 10^{14}$ per bunch. The peak SEY is assumed to be 1.8. The peak neutralization level is about 1.5% within the rms beam radius, and about 10% on average within the beam pipe.

(Figs. 21, 22). These electrons back-scatter from the catcher and the vacuum chamber, resulting in a high concentration of electrons with a broad energy spectrum. The injecting- and circulating-beams impacting on the foil produce secondary emission of electrons at tens of eV energy. The injecting- and circulating-beam also produce knock-on electrons at a high energy (up to several MeV). The stripping-foil, operating at a high temperature around 2000 K, emits thermionic electrons at low energy [36].

An accurate control of electron collection is essential in minimizing the electron cloud in the injection region. As shown in Fig. 23, the catcher consists of multiple tapered blades of low charge-state (carbon-carbon) to reduce back-scattering [38]. Video cameras are used to monitor both the stripping foil and the electron catcher (Fig. 24).

4.2 Collimation region electrons

Protons incident on the collimator surfaces produce secondary electrons. Depending on the energy of the beam and the incident angle, the secondary electron-to-proton yield can greatly exceed 1 when the incident beam is nearly parallel to the surface [39], as shown in Fig. 25. Although a serrated surface reduces the generation of secondary-emission electrons, it is practically ineffective since protons incident on the front edge of the teeth easily escape from the collimator body due to the long proton stopping-length (about one meter) at 1 GeV energy. It is crucial to use a two-stage collimation system with a primary thin scraper (Fig. 26) to enhance the impact parameter to the secondary collimator (Fig. 27).

The SNS ring collimation system is designed to collect...
Figure 23: Trajectory of the incident and back-scattered electron from the SNS injection electron catcher [38].

Figure 24: SNS injection foil chamber and nearby devices showing video ports to monitor the stripping foil and the electron catcher.

Figure 25: Proton-induced secondary-emission yields of electrons as functions of the incident angle for 28-MeV protons striking a flat (solid line) and a serrated (open circles) stainless-steel surface (courtesy P. Thieberger).

Figure 26: Schematics of the SNS ring primary scraper assembly and the down-stream shielding. The scraper assembly contains 4 thin tantalum blades transversely placed at 45° angle. The down-stream shielding has a similar structure as the ring collimators shown in Fig. 27 (courtesy H. Ludewig and N. Simos).

Figure 27: Schematics of the SNS ring collimator showing layers of material for radio-activation containment. The effective length is about 1.5 m. The collimator is designed to withstand an average beam power of up to 10 kW at 1 GeV kinetic energy (courtesy H. Ludewig, N. Simos, and G. Murdoch). The SNS ring has two such collimators, both located at optimized betatron phase advances down-stream of the primary scraper and its shielding [40].

more than 90% of the lost beam in the ring [40]. Comparing with other regions where the beam loss is typically below 1 W/m for hand-on maintenance, the collimation region expects a beam loss of about 500 W/m. In order to avoid frequent maintenance in such high radio-activated region, the secondary collimators are made non-adjustable (Fig. 27) without attaching any electron-cloud mitigation devices (like clearing electrodes and solenoids).

The small aperture of the collimators (Fig. 28) helps to reduce the trailing-edge electron multipacting in the most critical region where beam loss is expected to be the greatest. As shown in Fig. 29, the lower energy gain in the collimator region results in an effectively lower secondary yield. In areas of larger aperture between the collimators, which are more susceptible to the multipacting, solenoids are wound to suppress the electron cloud.
4.3 Ionization

The rate of electron line-density increase per unit length of circumference is given by the relation

\[
\frac{d^2 \lambda_e}{d \beta d s} = \frac{\rho_m \beta \sigma_{ion} P}{e}
\]

where \( I \) is the beam current, \( \sigma_{ion} \) is the cross-section, \( P \) is in units of Torr. At the room temperature of 300 K, the molecular density \( \rho_m \) is about \( 3.3 \times 10^{22} \text{ m}^{-3} \). For the SNS ring at a pressure of \( 10^{-8} \text{ Torr} \), a total number of \( 2.6 \times 10^9 \) electrons per turn, or \( 1.7 \text{ pC/m} \) of charge per turn, is produced when the proton intensity is \( 2 \times 10^{14} \). The ionization electrons, however, are likely to be trapped in the long proton bunch until the very end of the bunch passage, thus not participating in the trailing-edge multipacting process [25].

5 MITIGATION MEASURES

Control of the electron-cloud effects in both the RHIC and the SNS ring involves design and operational optimization to minimize the uncontrolled beam loss, enhancement of Landau damping, and suppression of electron generation.

Minimization of the uncontrolled beam loss is of primary importance for the SNS ring [6]. At the design stage, a large transverse acceptance and long, uninterrupted straight-sections are two important aspects considered on the lattice. The two-stage collimation system also plays an essential role in localizing the beam loss to controlled areas. For RHIC, intra-beam scattering causes routine beam loss during the beam storage. The collimation system can effectively control the vacuum pressure rise and reduce the experimental background.

Enhancement of Landau-damping helps to raise the instability threshold. For the SNS ring, a large vacuum-pipe aperture in the high-dispersion area and a large RF voltage provide sufficient momentum acceptance of \( \pm 1\% \) for a 480 \( \mu \text{m} \) transverse acceptance; longitudinal painting is used to expand the momentum spread of the injected beam; lattice sextupole families are used for chromatic adjustments, to either improve momentum acceptance or enhance damping; and, octupoles are also available for a possible future enhancement. For RHIC, octupole correction magnets are used to suppress the transverse instability near transition crossing.

Suppression of electron generation mainly includes measures to reduce the electron multipacting as discussed in the following.

5.1 Surface treatment

The inner surface of the stainless-steel pipe in the RHIC warm region are coated in subsequent sections with non-evaporative getter (NEG) to lower both the secondary emission and desorption yield, and to increase local vacuum pumping [29]. Beam experiments confirmed the effectiveness of the coating in reducing the vacuum pressure rise [42]. In year 2004, part of the RHIC cold region (four arc sections and four triplet sections) is pumped down to a vacuum pressure below \( 1 \times 10^{-2} \text{ Torr} \) before the cryogenic cool down to reduce the physically sorbed gas to submonolayers.

The inner surface of the entire SNS ring is coated with TiN. Coated chambers are tested for conductivity and risetime response. With a variation within 20\%, the thickness of 0.1 \( \mu \text{m} \) withstands the bombardment from electrons during the lifetime of the machine operation. Two layers of coating are applied to the ceramic chamber for injection kickers: a 1 \( \mu \text{m} \)-thick copper layer for by-passing the image charge, and a 0.1 \( \mu \text{m} \)-thick titanium nitride (TiN) layer for a low secondary-electron yield, along with an exterior
metal enclosure for dc current by-pass. Such a design allows the passage of the image current above a frequency of the lowest betatron sideband (~200 kHz) without degrading the magnetic-field penetration (a rise time of about \( 200 \mu s \)), eddy-current heating, and beam-induced heating. For the ferrite of the extraction kicker inside the vacuum pipe, a special coating pattern is selected to avoid eddy-current heating, rise-time degradation, and high-voltage shorts, as shown in Fig. 30. Measurements of TiN-coated surfaces indicate a reduction of SEY by more than half of an unit (Fig. 3) [27]. Fig. 31 shows the expected reduction of the electron density with a partial and a full TiN coating. Vacant ports are available for additional pumping, if needed, to accommodate a higher level of outgassing from the rougher coated surface. The present magnet and vacuum chamber design does not allow NEG film coating which requires in-situ baking.

Beam scrubbing at an electron dose level of 1 mC/mm\(^2\) can further condition the vacuum surface during operations [43]. For the SNS ring, vacant ports can house turbo pumps to function at a vacuum pressure above \( 10^{-6} \) Torr. Extended beam storage in the ring is planned to accelerate the scrubbing process.

5.2 Solenoids

Solenoidal magnetic field helps to confine the motion of electrons to be in parallel to the vacuum pipe surface, effectively reducing the electron multipacting. Fig. 32 shows the measured reduction of vacuum pressure in a RHIC warm section when the solenoidal fields are applied [14].

Solenoids are implemented on the vacant straights (about 5 m) of the collimation region of the SNS ring. The solenoid field \( B_0 \) of about 50 G is so strong that the radius \( r_0 = m_e v_t / e B_0 \) (12 mm for 300 eV electrons) of electron motion is small compared with the pipe radius, as shown in Fig. 33. Effects on the proton beam can be minimized by alternating the polarities of the solenoids observing betatron-phase variations. Skew quadrupoles can further be used to correct the coupling.
5.3 Clearing electrodes

Clearing electrodes have long been used to mitigate electron-cloud effects in a coasting beam by locally changing the beam potential that traps the electrons [2, 3]. For a bunched beam, the electrodes can alter the multipacting pattern (frequency and energy gain) to effectively suppress the electron generation even if the applied voltage is much less than the beam potential, as shown in Fig. 34 [44].

For the SNS ring, a dedicated clearing electrode is implemented inside the stripping-foil assembly (Fig. 22). A voltage up to 10 kV can be applied, adequate to suppress multipacting in the injection region. BPMs around the ring are designed to be also used as clearing electrodes capable of providing a voltage up to ±1 kV, adequate to locally reduce the multipacting frequency (Fig. 35).

5.4 Others

In order to reduce the electrons trapped by the residual beam in the beam gap of the SNS ring, a gap-clearing kicker is designed to clear the 250 ns gap during the last 100 of the 1060-turn accumulation [46]. By resonantly exciting coherent betatron oscillations, beam residuals are driven into the primary collimator, where beam loss is measured with a gated fast loss monitor. For the SNS ring, vacuum ports are screened, and steps in the vacuum pipe are tapered at 1-to-3 ratio to reduce peaked electric fields causing electron emission. A vacuum pressure of about $10^{-8}$ Torr ensures low electron generation from gas ionization. Possibilities of mitigating the e-p instability with wide-band resistive feedback are also under investigation [47].

6 BEAM-DYNAMICS CHALLENGES

Electron-cloud effects limit the RHIC upgrade path toward higher intensity and luminosity, and the SNS potential beyond the design beam intensity and power. During the past four decades, many progresses have been made in the field of electron cloud. However, the understanding of the subject remains incomplete.

Among many open challenges are (1) understanding of surface science underlying secondary emission and gas desorption and guidance on surface treatment (2) dynamic model on vacuum pressure rise incorporating desorption, pumping, and ionization, and comparison with measurements; (3) instability theory and simulation that satisfactorily reproduces long-bunch regime data (PSR) and predicts...
high-intensity ring performance (SNS, J-PARC); (4) self-
consistent treatment of electron-cloud formation and beam
instability in the simulation codes; (5) systematic bench-
marking of simulation codes, and benchmarking with ex-
perimental data; (6) reliable diagnostics for both warm and
cold regions; and (7) wide-band damping of fast, single-
bunch electron-cloud instability.

We are indebted to our colleagues in the e-cloud com-
munity, and to colleagues on RHIC and SNS projects, es-
specially A. Aleksandrov, G. Arduini, M. Bai, J. Brodowski,
P. Cameron, N. Catalan-Lasheras, A. Chao, R. Connolly,
S. Cousineau, V. Danilov, D. Davino, A. Fedotov, M.
Furman, O. Gröbner, H. Hahn, P. He, S. Henderson, N.
Hilleret, H. Huang, Y. Y. Lee, H. Ludewig, R. Macek,
W. Meng, R. Michnoff, C. Montag, M. Pivi, M. Plum, C.
Prior, V. Pitsin, D. Raparia, G. Rees, N. Simos, F. Ruggi-
tero, T. Satogata, H. Schonauer, N. Simos, S. Tepikian, P.
Thieberger, R. Todd, J. Tuozzolo, F. Zimmermann, B. Zot-
ter. We thank F. Zimmermann for the careful reading and
corrections of the manuscript.

7 REFERENCES
Electron and Positron Storage Rings, Saclay (1966, Orsay,
Univ. de France) VIII-6-1
(1971)
Accel. Conf., edited by P. Lucas, S. Webber (IEEE, Chicago,
2001), p. 688
S. Webber (IEEE, Chicago), p. 666
5044
Japan (1998)
Beam Dynamics Newsletter, 33 (2004) 128; BNL Tech. note
C-A/AP/191 (2005)
[15] See, for example, articles in the ICFA Beam Dynamics
3 (2000) 080101; J. Wei, M. Blaskiewicz, J. Brodowski, et al,
[17] F. Zimmermann, LHC Project-Report 95, and SLAC-PUB-
7425 (1997).
[23] L. Wang, A. Chao, H. Fukuma, ECLASS 2004 (Napa,
2004).
[27] P. He, H. Hseuh, M. Mapes, et al, submitted to J. Vacuum
[28] A. Gröbner, these proceedings
[29] H. Hseuh, et al, ECLASS 2004 (Napa, 2004); private com-
munications (2005)
[30] J. Wei, Longitudinal Dynamics of the Non-Adiabatic Regime
in the Alternating-Gradient Synchrotrons, Ph.D. dissertation
(1989, revised 1994, Stony Brook)
ST-AB, 5 (2002) 084401
[33] Beam experiments participated by J. Wei, U. Irimo, M. Bai,
M. Blaskiewicz, P. Cameron, R. Connolly, W. Fischer, H.
Huang, R. Lee, R. Michnoff, V. Pitsin, T. Roser, T. Satogata,
S. Tepikian, L. Wang, S.Y. Zhang, to be published (2005)
SNS/59 (1999)
Lett., 93 042901.
(2004) 034401
PRST-AB 4 (2001) 010101
Optics Design for a Nonlinear Collimation System in the LHC

A. Faus-Golfe, J. Resta López IFIC, Valencia, Spain
F. Zimmermann, CERN, Geneva, Switzerland

Abstract

We describe the adaptation of a nonlinear collimation system previously developed for CLIC for LHC betatron cleaning. The LHC nonlinear system employs 1 or 2 pairs of skew sextupoles, e.g., one for the horizontal and one for the vertical collimation. We discuss the optics of this system, and outline a possible plan for further work.

1 INTRODUCTION

A collimation system for the LHC should (1) prevent beam-loss induced quenches if the superconducting LHC magnets; (2) minimize activation of accelerator components outside of the dedicated collimation insertions; (3) ensure an acceptable background in the particle-physics experiments; (4) withstand the impact of eight bunches in case of an irregular beam dump; and (5) not introduce intolerable wake fields that might compromise beam stability [1]. These are similar requirements as those which were posed for the nonlinear collimation design of CLIC [2]. It is thus a close thought to apply the same design scheme to the LHC.

The main differences from CLIC are the following:

- the LHC momentum spread is almost two orders of magnitude smaller, and, hence, cannot be exploited for widening the beam during collimation;
- emittance growth from synchrotron radiation is insignificant, and does not constrain the design of the collimation system;
- the geometric vertical emittance is about 3 orders of magnitude larger than in CLIC.

2 SCHEME

The basic layout of a nonlinear collimation system is illustrated in Fig. 1. The purpose of the first nonlinear element is to blow-up beam sizes and particle amplitudes, so that the collimator jaw can be placed further away from the nominal beam orbit.

At each nonlinear element a particle suffers deflections \( \Delta q'_i = -\partial H_n / \partial q_i \), where \( H_n \) is the Hamiltonian of the multipole. As was pointed in Ref. [3], higher-order multipoles (decapoles, dodecapoles, etc.) are not useful, because they do not penetrate to the small distances necessary. Skew sextupoles and octupoles could be used.

Figure 1: Schematic of a nonlinear collimation system.

3 HISTORY

Some types of nonlinear collimation systems for future linear colliders have been described in the literature [2, 3, 4, 5]:

- For the NLC, Merminga et al. [3] proposed a scheme with skew-sextupole pairs for nonlinear collimation in the vertical plane.
- Subsequently, R. Brinkmann, R. Pitthan, P. Raimondi, A. Seryi et al. [4], presented a halo reduction method with the addition of “tail-folding” octupoles (“Chebyshev arrangement of octupoles”) in the NLC final focus system (see also [6] for an earlier study with only 1 octupole in front of the final doublet).
- For the TESLA post-linac collimation system a magnetic energy spoiler (MES) has been suggested [5]. A octupole is placed at a high dispersion point between a pair of skew sextupoles (at \( \pi / 2 \) phase advance from the octupole). The skew sextupoles are separated by a optical transfer matrix \(-I\). The result is a significant increase in the vertical beam size at a downstream momentum spoiler.

A characteristic feature of all these systems is that they separate between energy and betatron collimation, and typically employ the nonlinear elements only in one or the other half.

A nonlinear collimation system for CLIC with a pair of skew sextupoles is being explored [2]. It presents a single vertical spoiler which collimates in the transversal betatron degrees of freedom and in energy. The scheme is illustrated in Fig. 2. More details of this system can be found in Ref. [2].
4 SYSTEM EQUATIONS FOR LHC

The Hamiltonian of a skew sextupole at a location with horizontal dispersion \( D \) is
\[
H_s = \frac{1}{6} K_s (y^3 - 3(x + D\delta)y) ,
\]
where \( x \) and \( y \) are the transverse betatron amplitudes at the sextupole, and \( \delta \) the relative momentum offset. The integrated sextupole strength \( K_s \) can be expressed in terms of the sextupole length \( l_s \), the pole-tip field \( B_T \), the magnetic rigidity \( (B\rho) \), and sextupole aperture \( a_s \) as
\[
K_s = \frac{2B_T l_s}{(B\rho)a_s^2} .
\]

At the skew sextupole a particle suffers deflections \( \Delta x' = -\partial H_s/\partial x, \Delta y' = -\partial H_s/\partial y \), or
\[
\Delta x' = K_s (D_{\text{sext}}\delta + x)y \\
\approx K_s x y
\]
\[
\Delta y' = -\frac{1}{2}K_s (y^2 - x^2 - D_{\text{sext}}^2\delta^2 - 2D_{\text{sext}}\delta x) \\
\approx -\frac{1}{2}K_s (y^2 - x^2) ,
\]
where in the second step we have neglected the dispersive term, since dispersion and energy spread are small.

The position at a downstream spoiler is obtained from
\[
x_{sp} = x_{0,sp} + R_{12}\Delta x' ,
\]
\[
y_{sp} = y_{0,sp} + R_{34}\Delta y' ,
\]
where the subindex \( 0 \) indicates the position in the absence of the skew sextupole and \( R_{12}, R_{34} \) are the optical transport matrix elements between the sextupole and the spoiler.

The rms beam size at the spoiler is computed by squaring the expressions for \( x_{sp} \) and \( y_{sp} \), and averaging over the transverse distribution:
\[
\sigma_x \approx (K_s^2 R_{12}^2 B_{x,\text{sext}} \beta_{y,\text{sext}} \epsilon_x \epsilon_y + \beta_{x,\text{sp}} \epsilon_x)^{1/2}
\]
\[
\sigma_y \approx (K_s^2 R_{34}^2 (3\beta_{x,\text{sext}}^2 \epsilon_x^2 + 3\beta_{y,\text{sext}}^2 \epsilon_y^2) \\
-2\beta_{x,\text{sext}} \beta_{y,\text{sext}} \epsilon_x \epsilon_y + \beta_{x,\text{sp}} \epsilon_x)^{1/2} .
\]

For spoiler survival in case of beam impact, a minimum beam size \( \sigma_{r,\text{min}} \) of about 200 \( \mu m \) is required [7], so that
\[
s_y \sigma_x \geq \sigma_{r,\text{min}}^2 ,
\]
which, using \( \epsilon \equiv \epsilon_x = \epsilon_y \), we can rewrite as
\[
(K_s^2 R_{12}^2 B_{x,\text{sext}} \beta_{y,\text{sext}} \epsilon_x + \beta_{x,\text{sp}} \epsilon_x) \times
\]
\[
[K_s^2 R_{34}^2 (3\beta_{x,\text{sext}}^2 + 3\beta_{y,\text{sext}}^2 \epsilon_y) \\
-2\beta_{x,\text{sext}} \beta_{y,\text{sext}} \epsilon_x \epsilon_y + \beta_{x,\text{sp}} \epsilon_x]^2) \geq \sigma_{r,\text{min}}^4 .
\]

This determines the minimum values of \( K_s, R_{12} \) and \( R_{34} \) required.

We denote the collimation amplitude defined by the skew sextupole for the horizontal and vertical betatron motion as \( \pm n_x \sigma_{x,\text{sext}} \) and \( \pm n_y \sigma_{y,\text{sext}} \), respectively, and the physical aperture of the spoiler by \( \pm n_x \sigma_{x,\text{sp}} \) and \( \pm n_y \sigma_{y,\text{sp}} \).

For the collimation to function in either transverse plane, we must have
\[
\frac{1}{2}K_s R_{12} n_x \sqrt{\beta_{x,\text{sext}} \sigma_{x,\text{sext}}^2 \epsilon_x} = n_{y,\text{sp}} \sqrt{\beta_{y,\text{sp}} \epsilon_y} 
\]
\[
\frac{1}{2}K_s R_{34} n_y \sqrt{\beta_{y,\text{sext}} \sigma_{y,\text{sext}}^2 \epsilon_y} = n_{x,\text{sp}} \sqrt{\beta_{x,\text{sp}} \epsilon_x} .
\]

Particles with \( |y_{\text{sext}}| \approx |x_{\text{sext}}| \) will not be collimated by the vertical spoiler. To catch these particles as well, we use the horizontal deflection by the skew sextupole and the horizontal aperture of the spoiler, \( n_x \sigma_{x,\text{sp}} \). We require
\[
K_s R_{12} n_x \sqrt{\beta_{x,\text{sext}} \sigma_{x,\text{sext}}^2 \epsilon_x} = n_{y,\text{sp}} \sqrt{\beta_{y,\text{sp}} \epsilon_y} 
\]
where the horizontal amplitude aperture at the spoiler \( n_x \) can be adjusted to improve the cleaning efficiency for particles with offsets in both transverse planes. For example, we might set \( n_x = 2n_{y,\text{sp}} \), if the horizontal and vertical beta functions are equal at the skew sextupole and we want to approximate a circular collimation aperture in the transverse normalized \( x - y \) space.

Now assuming \( \beta_{x,\text{sext}} = \beta_{y,\text{sext}} \), from Eq. (13) we can obtain the beta function at the skew sextupole as
\[
\beta_{x,\text{sext}} = \left( \frac{n_{x,\text{sp}}^2 \beta_{x,\text{sp}}}{K_s^2 R_{12}^2 n_y^2} \right)^{1/2} .
\]

In this case, we may rewrite the condition (9) for spoiler survival in terms of collimation amplitudes as
\[
\frac{R_{12}^4 n_x^2}{R_{12}^2 n_x^2} \left( \frac{n_{y,\text{sp}}^2 + 1}{n_{x,\text{sp}}^2 + 1} \right) \beta_{x,\text{sp}}^2 \geq \sigma_{r,\text{min}}^4 .
\]

Fig. 3 shows the product \( \sigma_y \sigma_x \) in units of \( \sigma_{r,\text{min}}^2 \) as function of \( n_x \) and \( n_{y,\text{sp}} \).

Unlike for the CLIC case, where off-momentum particles were collimated simultaneously, we here still have some freedom in selecting the parameters. The tightest constraint likely arises from the achievable skew sextupole strength.

5 OPTICS

Here we limit our study to the optics for a system composed of a single spoiler and a pair of skew sextupole as
illustrated in Fig. 4. The first skew sextupole blows-up the particle amplitudes, thereby allowing larger collimator gaps which may avoid unacceptable high transverse resistive impedance from the collimator material. A skew sextupole downstream of the spoiler, and at \( \pi \) phase advance from the first sextupole, cancels the geometric aberrations induced by the former.

Let us take as an example \( n_x = n_y \approx 6 \) and \( n_{y^2} = 8 \). Several nonlinear-collimation optics for IR7 in LHC optics version V6.5 have been matched, 12 in total. All of these optics solutions fulfil the above equations and requirements. They differ in the beta functions at the skew sextupoles and in the \( R_{12}, R_{34} \) matrix elements between the first skew sextupole and the spoiler. The matching was done without affecting the optics of the other LHC insertions, and involved only existing quadrupole magnets. To elucidate which of the different optics solutions is best suited for our application, we can choose a number of criteria:

- minimize the normalized sextupole strength \( K_s \),
- minimize the product of normalized sextupole strength and the larger of the two beta functions at the skew sextupole (which is equivalent to minimizing the sextupole pole-tip field),
- minimize the nonlinear aberration introduced by the first skew sextupole, which scale as \( \beta_{y,\text{sext}} \beta_{x,\text{sext}}^{3/2} K_s \) and as \( \beta_{1,\text{sext}}^{1/2} \beta_{y,\text{sext}}^{1/2} K_s \).

From the 12 candidate solutions, we have selected an optics for which both the sextupole strength and the product of sextupole strength and beta function at the sextupole are minimum. Table 1 summarizes the main parameters of this optics including the strength of the skew sextupoles. The corresponding plot for the beta functions is given in Figure 5. It can be compared with the present LHC optics of IR7 V6.5 in Fig. 6.

Figure 3: Surface \( \sigma_x \sigma_y / \sigma_{r,\text{min}}^2 \) as function of \( n_x \) and \( n_{x^2} \). The point represents the working point \( n_x = 6, n_{x^2} = 2n_y = 16 \) and \( \sigma_x \sigma_y / \sigma_{r,\text{min}}^2 = 1.425 \). The solid line in the plane \( n_{x^2} \) vs. \( n_x \) represents the limit \( \sigma_x \sigma_y / \sigma_{r,\text{min}}^2 = 1 \) for spoiler survival.

Figure 4: Schematic of a nonlinear collimation layout for the LHC.

Figure 5: The optics solution proposed for LHC IR7 with a nonlinear collimation section based on two skew sextupoles.

Figure 6: The present optics of LHC IR7 (V6.5).
Table 1: Optics parameters for a nonlinear collimation section in IR7 of LHC.

<table>
<thead>
<tr>
<th>variable</th>
<th>value</th>
</tr>
</thead>
<tbody>
<tr>
<td>beta functions ((x, y)) at skew sext.</td>
<td>200.0, 200.0 m</td>
</tr>
<tr>
<td>product of skew sextupole pole-tip</td>
<td></td>
</tr>
<tr>
<td>field and length ((B T f_s))</td>
<td>8.1823 T·m</td>
</tr>
<tr>
<td>skew sextupole aperture (a_s)</td>
<td>10 mm</td>
</tr>
<tr>
<td>skew sextupole strength (K_s)</td>
<td>7.0063 m⁻²</td>
</tr>
<tr>
<td>(R_{12}, R_{34}) from sext. to spoiler</td>
<td>124.403, 124.404 m</td>
</tr>
<tr>
<td>beta functions ((x, y)) at spoiler</td>
<td>77.381, 77.381 m</td>
</tr>
<tr>
<td>rms spot size ((x, y)) at spoiler</td>
<td>215.89, 263.96 µm</td>
</tr>
</tbody>
</table>

**6 COLLIMATION APERTURE**

From Eqs. (3) and (4), and knowing the optics parameters at the sextupoles and the spoiler, we can compute the collimation contours, defined by the deflections at the first skew sextupole, as

\[
ny^2 \frac{\beta_{y, \text{aperture}}}{\beta_{x, \text{sext}}} = \frac{1}{2} K_s R_{34}(X^2 - Y^2), \quad (16)
\]

\[
nx^2 \frac{\beta_{x, \text{aperture}}}{\beta_{x, \text{sext}}} = K_s R_{12} XY, \quad (17)
\]

where we use horizontal and vertical coordinates normalized with the transformations \(X = x/\sqrt{\beta_x}\) and \(Y = y/\sqrt{\beta_y}\). Figure 7 shows the resulting collimation boundaries. Note that the boundaries shown refer to vanishing initial slopes, and that they would be modified for trajectories with initial \(x'\) or \(y'\) unequal zero.

![Collimation contours](image)

**7 SUMMARY AND OUTLOOK**

Following the principles of nonlinear collimation, previously studied for TeV linear colliders, we propose the adaptation of a nonlinear collimation system for the LHC. It employs 1 pair of skew sextupoles separated by a \(-I\) transfer matrix in order to cancel geometric aberrations. The deflection from the first sextupole defines the collimation amplitude for the horizontal and vertical betatron motion, which is related to the aperture of the spoiler by the Eqs. (11), (12) and (13). A scheme with 2 pairs of skew sextupoles might be studied with similar features.

Choosing \(\beta_{x, \text{sext}} = \beta_{y, \text{sext}}\) and assuming basic requirements and criteria, such as the minimum beam size \((\sigma_x, \sigma_y \approx 200 \mu m)\) for spoiler survival in case of beam impact, we have calculated example optics parameters. The tightest constraint likely arises from the achievable skew sextupole strength.

Here our purpose has been to demonstrate the principal viability of nonlinear collimation for the LHC. Further works should add secondary absorbers (at a phase advance \(\mu_0\) close to the optimum \(\pm \arccos(n_y/n_y)\) [8], or about 41° and 131°, behind the spoiler) and compare the cleaning efficiency of the full system with that of a conventional linear collimation system by means of multi-particle tracking.

**8 REFERENCES**

[6] K. Thompson, R. Pitthan, F. Zimmermann, et al., NLC Collimation Meetings, in particular 22.05.98, 29.05.98, and 31.08.98.; see web site: http://www-project.slac.stanford.edu/lc/bdir/meetings/collimation.asp
What is an acceptable dynamic vacuum pressure in the LHC arcs?

B. Jeanneret, F. Zimmermann, CERN, Geneva, Switzerland

Abstract

Beam-gas nuclear and Coulomb interactions introduce limitations on the tolerable vacuum pressure in the LHC and its upgrade. We discuss local pressure limits set by the heat load on the cryogenics system and by loss-induced quenches, as well as average pressure limits imposed by beam-lifetime degradation and emittance blow up. These pressure limits depend on the molecular species in the residual gas. We also consider the effect of Coulomb scattering, which strongly depends on the beam energy. An important ingredient to our discussion is the available cooling capacity, which we estimate for the nominal and ultimate LHC, as well as for two possible upgrade scenarios. The quoted pressure tolerances refer to the nominal LHC only. They would become significantly tighter for higher beam currents or for shorter bunches.

1 INTRODUCTION

Residual-gas molecules reduce the proton beam lifetime via nuclear scattering and they increase the beam emittance by multiple elastic Coulomb scattering. The minimum sustainable beam lifetime is determined by the associated heat load on the cold bore and the available cooling capacity. The magnitude of local pressure bumps is limited by magnetic quenches. The acceptable average gas pressure is also constrained by the emittance growth on the LHC injection plateau (about 20 minutes), which should not exceed a few percent.

In the following sections, we first review the available cooling capacity at two temperatures accounting for some expected basic heat loads. We next infer limits on the average and local vacuum pressure, imposed by the heat load from beam-gas scattering, together with the available cooling capacity and expected quench limits. We then discuss nuclear scattering processes and the pertinent loss distributions in a little more detail. In the following, the emittance growth from multiple Coulomb scattering and single Rutherford scattering are addressed. Finally, we draw some conclusions.

2 COOLING CAPACITY

Table 1 lists estimated head loads at 4.6–20 K from image currents, static heat inleaks, and synchrotron radiation at injection and top energy, as well as the available cooling capacity, for the nominal and ultimate LHC and also for two possible LHC upgrade scenarios [3]. The installed cryogenics capacity for the 4.6–20 K temperature level is about 24 kW per sector [4]. However, the cooling capacity at 4.6–20 K is limited by the hydraulic impedance of the beam-screen cooling loop to about 2.4 W/m per aperture [4] as listed in Table 1. Without the hydraulic impedance limit, the installed refrigeration capacity would correspond to a significantly larger value of 3.9 or 4.2 W/m per aperture for the upgraded old and the new cryoplants, respectively [4].

The LHC arc magnets feature a ‘double aperture’ for the two separated beams. Throughout this paper, we quote the cooling capacity and heat load per aperture, i.e., for one beam pipe. The values per meter of magnet length would be two times larger.

The aforementioned capacities should be compared with the heat loads at 4.6–20 K in Table 1. Subtracting the unavoidable heat loads from static inleaks, synchrotron radiation, and image currents, for the nominal LHC a capacity of about 2 W/m remains to remove heat induced by electron cloud or by ions bombarding the surface. Assuming standard scaling laws, the available capacity decreases to 1.7 W/m, for the ultimate LHC, to about 0 W/m for the baseline upgrade scheme with more and shorter bunches, and to a more acceptable 1 W/m for the Piwinski-type upgrade with fewer longer bunches.

At 1.9 K the cooling capacity is about 0.5 W/m. The sum of nominal static and dynamic load per arc cell is 40.9 W, which corresponds to 0.19 W/m per aperture. This is more than two times smaller than the estimated cooling capacity, which may suggest a significant margin. Also note that cooling 0.2 W/m at 1.9 K is equivalent to cooling 1 W/m at 4.6–20 K.

3 BEAM LIFETIME AND HEAT LOAD

For the first time in an accelerator, at the LHC the beam loss due to nuclear scattering represents a non-negligible heat load on the cold bore of the magnets. In consequence, a minimum beam lifetime of 100 h has been required, so as to ensure that the heat load due to nuclear scattering stays below 0.1 W/m for the two beams or below 0.05 W/m per aperture [5].

The nuclear cross section consists of two parts: one for inelastic and one for elastic scattering. As discussed further below, inelastic interactions lead to a local loss, while...
elastic or quasi-elastic scattering events may lead to losses at the collimators and do not necessarily contribute to the local losses. At high energies, the elastic part of the nuclear interactions amounts to about 20% of the total cross section for hydrogen. This fraction increases for heavier atoms, but it always stays below 50% (optical theorem). Beam-gas scattering at a lifetime of 100 h translates into an average heat load at 1.9 K between 0.024 W/m and 0.038 W/m per aperture, depending on whether or not elastic scattering (here taken to amount to 40% of the total cross section) leads to a local loss.

The beam lifetime $\tau$ is related to the total nuclear interaction cross section $\sigma_{\text{total}}$ and gas density $n$ via $1/\tau = c \sigma_{\text{total}}$, where $c$ denotes the speed of light. Table 2 lists the total and inelastic nuclear-interaction cross sections for various gas species. Most of these cross sections were taken from [6]. Analytical formulae for the inelastic cross section can be found in [7, 8]. The next column in Table 2 contains the maximum allowed molecule densities $n$ inferred from the 100-h lifetime condition. The highest density of $1.0 \times 10^{15} \text{ m}^{-3}$ is permitted for hydrogen. It corresponds to a pressure of 32 ntorr at room temperature.

During commissioning the beam lifetime could be shorter than the 100-h design value and the heat load larger accordingly. For a beam lifetime of 1 h at 7 TeV and nominal beam intensity, the heat load at 1.9 K would be 2.4–3.8 W/m, which is unacceptable. Since the 0.5-W/m cooling capacity at 1.9 K has a margin of about 0.3 W/m, we could accept a beam lifetime up to about 10 times shorter than 100 h, i.e., 10 h. This is about the lowest lifetime from nuclear interactions that could be achieved with the nominal beam intensity. This limit assumes that the gas pressure is uniform around the ring, and that there is zero heat-load contribution from the electron cloud (on the other hand it also does not make use of any possibly available spare cooling capacity at 4.6–20 K). The foregoing analysis is based on a steady-state analysis, which appears appropriate for the residual-gas pressure. The analysis does not cover transient conditions, e.g., drops of the beam lifetime to 10 minutes for a short period of time.

During scrubbing, electron-cloud heat load and vacuum pressure are most likely correlated. Then part of the 1.9-K cooling capacity may be needed to cool the beam screen at 4.6–20 K, heated by the electron cloud, or vice versa, and in this case the minimum acceptable lifetime for the nominal beam could be modified.

### 4 QUENCHES AND LOCAL PRESSURE

The dipole magnet quench limit at 7 TeV corresponds to a local loss rate of $8 \times 10^6 \text{ p/m/s}$, and that at 450 GeV to a loss rate of $7 \times 10^6 \text{ p/m/s}$. For comparison, a 100 h beam lifetime corresponds to $6 \times 10^4 \text{ p/m/s}$ adding the losses from the two beams. Therefore, at 7 TeV the dipole magnet quench limit is reached for a gas density about two orders of magnitude higher than the average loss rate for a 100-h lifetime (see also Ref. [9], p. 70).

### 5 ELASTIC INTERACTIONS

A fraction of the nuclear interactions are elastic or quasi-elastic (diffractive). This elastic or quasi-elastic component rises from 15% for hydrogen to almost 40% for heavier atoms. For atoms consisting of several nucleons, the diffractive interactions themselves are composed of two parts, namely diffraction on the atomic core as a whole and diffraction on individual nucleons inside the nucleus.

The simpler case are proton-hydrogen interactions. Introducing the Lorentz invariant $t$ parameter, $t \approx (p \theta)^2$ (with $p$ the momentum and $\theta$ the scattering angle), the differential cross section of purely elastic scattering can be expressed, approximately, as [10]

$$\frac{dt}{d\sigma} = \alpha e^{-b(t,s)t}, \quad (1)$$

where $\alpha$ represents a constant coefficient and $b(t, s)$ is the slope factor of the elastic differential cross section. If the
Table 2: Nuclear scattering total and inelastic cross sections, the implied maximum allowed residual-gas densities, and the accompanying emittance growth at injection for various gas species.

<table>
<thead>
<tr>
<th>molecule</th>
<th>$\sigma_{\text{tot}}$ [barn] at 7 TeV</th>
<th>$\sigma_{\text{el}}$ [barn] at 7 TeV</th>
<th>$n$ [m$^{-1}$] at $\tau_{\text{nucl}} = 100$ h</th>
<th>equiv. pressure at 300 K [atm]</th>
<th>rad. length $X_0$ [g/cm$^2$]</th>
<th>$\tau_\epsilon$ [h] at 450 GeV for $\tau_{\text{nucl}} = 100$ h</th>
</tr>
</thead>
<tbody>
<tr>
<td>$H_2$</td>
<td>0.09</td>
<td>0.076</td>
<td>$1.0 \times 10^{15}$</td>
<td>32.0</td>
<td>61.3</td>
<td>28.4</td>
</tr>
<tr>
<td>$He$</td>
<td>0.133</td>
<td>0.102</td>
<td>$7.0 \times 10^{14}$</td>
<td>21.7</td>
<td>94.3</td>
<td>32.3</td>
</tr>
<tr>
<td>$CH_4$</td>
<td>0.511</td>
<td>0.383</td>
<td>$1.8 \times 10^{14}$</td>
<td>5.6</td>
<td>46.2</td>
<td>15.2</td>
</tr>
<tr>
<td>$H_2O$</td>
<td>0.510</td>
<td>0.368</td>
<td>$1.8 \times 10^{14}$</td>
<td>5.7</td>
<td>38.1</td>
<td>19.0</td>
</tr>
<tr>
<td>$CO$</td>
<td>0.751</td>
<td>0.523</td>
<td>$1.2 \times 10^{14}$</td>
<td>3.8</td>
<td>54.4</td>
<td>15.0</td>
</tr>
<tr>
<td>$CO_2$</td>
<td>1.171</td>
<td>0.754</td>
<td>$7.9 \times 10^{12}$</td>
<td>2.5</td>
<td>36.2</td>
<td>9.9</td>
</tr>
</tbody>
</table>

$t$-range of interest corresponds to scattering angles of order 10 times the LHC rms beam divergence, the $t$ dependence of $b$ may be neglected, and $b$ can be considered as a function of the centre-of-mass energy $\sqrt{s}$ alone [10].

For LHC energies one finds $b \approx 14 \ (\text{GeV}/c)^{-2}$ (see Fig. 3.3 in [10]), and the rms angular spread resulting from elastic proton-proton scattering is

$$\theta_{\text{rms}} \approx \frac{1}{\sqrt{b_p}} \approx 5.9 \times 10^{-4} \ \text{rad}.$$  (2)

This should be compared with the typical rms beam divergence, which at 450 GeV is estimated as

$$\theta_{\text{rms}} \approx \frac{1}{\sqrt{b_{pN}}} \approx 2.5 \times 10^{-4} \ \text{rad}.$$  (3)

Since the rms scattering angle is close to 100 times larger than the typical rms beam divergence at injection, and the aperture of the arcs is only of the order of 10–15 times the rms beam size, elastic proton-proton collisions almost exclusively lead to local losses.

Single diffractive scattering events for hydrogen have a cross section similar to the elastic one given above [10]. The $b$ parameter is also of the same order and, therefore, the same argument as for the elastic interactions applies, namely, protons diffracted off hydrogen nuclei are lost locally. At top energy, on the other hand, a significant part of the elastic interactions may lead to a particle loss at the collimators.

For more complex nuclei consisting of several nucleons, like oxygen, there are three elastic or quasi-elastic processes [10]: elastic scattering with one of the $n_{\text{eff}}$ nucleons inside the nucleus, single elastic diffraction between the incoming proton and a nucleus weighted with an effective nucleon number $n_{\text{eff}}$, and elastic coherent scattering off the entire nucleus. For the first two processes, the effective number of nucleons is approximately given by [10]

$$n_{\text{eff}} \approx 1.6 A^{1/3}.$$  (4)

The fraction of the total cross section corresponding to these events decreases for larger values of $A$, since the inelastic cross section grows more rapidly, as $A^{0.71}$. The parameter $b$ is roughly the same as for proton-hydrogen scattering. Hence, again, at injection energy, elastically scattered or singly diffracted protons are lost not far downstream of the scattering event.

A more important elastic scattering process for complex nuclei is proton-nucleus elastic scattering. The scattering distribution is still given by an equation of the form (1), where the parameter $b$ for this process now is [10]

$$b_{pN} \approx 14.1 A^{2/3} \ \text{GeV}^{-2},$$  (5)

which yields $b \approx 80 \text{ GeV}$ for $O, C, \text{ and } N$. At 450 GeV, the rms scattering angle for the elastic proton-nucleus scattering is

$$\theta_{\text{rms}} \approx \frac{1}{\sqrt{b_{pN}^p}} \approx 2.5 \times 10^{-4} \ \text{rad}.$$  (6)

which is about 30 times larger than the rms beam divergence, so that local losses are dominant. At 7 TeV, the rms scattering angle is $1.6 \times 10^{-5} \ \text{rad}$, equal to about 8 times the rms beam divergence. Here, roughly half of the elastically scattered protons will be intercepted by a collimator, while the others may circulate for a longer time in the ring.

Therefore, when computing either local losses or beam lifetimes, at injection energy the total nuclear cross section should be considered rather than the inelastic cross section alone. On the other hand, at top energy the total cross section determines the beam lifetime only, whereas the inelastic cross section characterizes the local losses.

6 EMITTANCE GROWTH

Multiple or single Coulomb scattering off the residual-gas atoms (mainly their nuclei) as well as elastic nuclear interactions can both increase the beam emittance.

The rms emittance growth due to multiple Coulomb scattering is [6, 11]

$$d(\gamma\epsilon) = \frac{\beta\gamma}{2} \left(\frac{136 \text{ MeV}}{\beta c_p}\right)^2 \frac{m_{\text{gas}}}{X_0} \frac{c n_{\text{gas}}}{X_0}.$$  (7)

were $m_{\text{gas}}$ is the mass of a molecule, $X_0$ the radiation length in units of kg m$^{-3}$, and $\beta$ the average beta function ($\bar{\beta} \approx 100 \text{ m at injection}$, and $\bar{\beta} \approx 150 \text{ m at top energy}$). The effect of multiple scattering is largest at injection. Emittance growth times expected at 450 GeV for
The partial gas pressures corresponding to a 100 h nuclear-scattering lifetime are also included in Table 2. The emittance doubling times are seen to range from 10 to 30 h, depending on the gas species. After some conditioning of the vacuum system, a fractional gas composition with less than 5% CO and CO₂ molecules is expected. The emittance growth times will therefore be dominated by H₂ or He. The multiple scattering formula (7) is valid to better than 11%, if the traversed distance in units of radiation length is larger than 10⁻³ (and smaller than 100). This gives rise to a lower time limit of validity for the multiple-scattering calculation which is of the order of 1 h.

7 SINGLE COULOMB SCATTERING

The integrated cross section for Mott scattering (the relativistic extension of Rutherford scattering) for scattering angles larger than a minimum angle θₘᵢₙ is [12]

\[ \sigma_{\text{Mott}}(θ > θ_{\text{min}}) = πα^2(\hbar c)^2 \frac{Z^2}{E_{\text{beam}}} \left[ \frac{1}{1 - \cos θ_{\text{min}}} - \cos θ_{\text{min}} - \ln(1 - \cos θ_{\text{min}}) \right]. \]

Tables 3 and 4 compare cross sections for single Coulomb scattering at angles above 1 or 10 times the rms beam divergence with the total nuclear-interaction cross section at two different beam energies, respectively. At injection, for the heavier atoms like oxygen or nitrogen Coulomb scattering at an angle larger than the beam divergence is about as probable as a nuclear interaction. However, in all cases the probability of an actual proton loss due to Rutherford scattering amounts to 1%, or less, of the loss probability from nuclear interactions.

8 CONCLUSIONS

We have estimated the cooling capacity available for electron cloud and beam-gas scattering at 4.6–20 K and at 1.9 K, respectively. From the latter we have derived limits on the acceptable beam lifetime and heat load, which translate into limits on the acceptable vacuum pressure. Additional tolerances on the vacuum pressure are imposed by magnet quenches due to local beam losses and by the emittance growth due to multiple Coulomb scattering off the residual gas.

In the nominal LHC at top energy, the spare cooling capacity per aperture is about 1.95 W/m for the beam screen, and 0.2–0.3 W/m for the cold bore. In the ultimate LHC, 1.7 W/m remain for beam-screen cooling, which shrinks to basically zero for the baseline upgrade and to 1.0 W/m for the alternative upgrade operating with fewer longer bunches, i.e., the so-called ‘Piwinski scheme’.

The minimum beam lifetime due to gas scattering sustainable by the LHC cryogenics system is of order 10 h at top energy. This would correspond to an average helium pressure of about 200 ntorr at 300 K. At injection energy, elastic nuclear scattering events lead to local losses, while at top energy elastically scattered protons are intercepted by the collimators. Inelastic scattering events — typically more than 60% of all nuclear interactions — always result in local losses. Single Coulomb scattering contributes at most to a 1% effect in the beam lifetime.

For a 100-h beam lifetime at injection, limited by nuclear scattering off the residual gas, the emittance growth due to multiple Coulomb scattering is of order 10–32 h. It varies with the gas species (32 h for helium, 10 h for carbon dioxide).

A local helium pressure of 2 μtorr, at 300 K, or alternatively, a local carbon dioxide pressure of 300 ntorr can induce magnet quenches in the arcs. These pressures correspond to an increase by a factor 200 above the design average value.

In this report, possible interplays of an elevated vacuum pressure with the electron cloud [13] were not discussed. The most direct interplay arises from the electron-induced vacuum-pressure increase, which leads to enhanced beam-gas scattering and raises the heat load on the cold bore of the magnets, while the electron cloud itself deposits heat on the beam screen. While the power deposition occurs on two different helium circuits, both are fed by a com-

<table>
<thead>
<tr>
<th>atom</th>
<th>( \sigma_{\text{tot}} ) [mbarn]</th>
<th>( \sigma_{\text{Mott}}(\theta &gt; \theta_{\text{min}}^{\text{beam}}) ) at 450 GeV/c [mbarn]</th>
<th>( \sigma_{\text{Mott}}(\theta &gt; \theta_{\text{min}}^{\text{beam}}) ) at 7 TeV/c [mbarn]</th>
</tr>
</thead>
<tbody>
<tr>
<td>H</td>
<td>45</td>
<td>0.07</td>
<td>0.46</td>
</tr>
<tr>
<td>He</td>
<td>133</td>
<td>0.29</td>
<td>1.32</td>
</tr>
<tr>
<td>C</td>
<td>331</td>
<td>3.6</td>
<td>16.5</td>
</tr>
<tr>
<td>N</td>
<td>379</td>
<td>3.5</td>
<td>22.5</td>
</tr>
<tr>
<td>O</td>
<td>420</td>
<td>4.6</td>
<td>29</td>
</tr>
</tbody>
</table>

<table>
<thead>
<tr>
<th>atom</th>
<th>( \sigma_{\text{Mott}}(\theta &gt; \theta_{\text{min}}^{\text{beam}}) ) at 7 TeV/c [mbarn]</th>
</tr>
</thead>
<tbody>
<tr>
<td>H</td>
<td>0.005</td>
</tr>
<tr>
<td>He</td>
<td>0.02</td>
</tr>
<tr>
<td>C</td>
<td>0.17</td>
</tr>
<tr>
<td>N</td>
<td>0.22</td>
</tr>
<tr>
<td>O</td>
<td>0.29</td>
</tr>
</tbody>
</table>
mon source and must, therefore, share the available cooling power. Scrubbing scenarios for the LHC must include the combined effect. The scrubbing strategy should aim to balance the distribution of cooling between cold bore and beam screen, taking into account the different conditioning behavior of molecular desorption and secondary emission [14].

9 ACKNOWLEDGEMENTS

We thank L. Tavian for several helpful and clarifying discussions.

10 REFERENCES

Minimum bunch length at the LHC

E. Vogel
CERN, Geneva
February 2005

Abstract

A luminosity upgrade of the LHC may be based on stronger focusing quadrupole magnets and/or a stronger focusing beam optics layout at the interaction regions, reducing the $\beta$-function from 55 cm to 25 cm. Due to the hourglass effect and the effects of the beam-beam interaction caused by the crossing angle, the bunch length becomes an important quantity for the luminosity. Within this scheme the reduction of the bunch length from nominal 0.31 m to 0.16 m leads to 43% more luminosity [1].

The paper presents estimates of the minimum bunch lengths at which Landau Damping is lost for three scenarios: (i) without additional measures, (ii) applying RF amplitude modulation and (iii) using a higher harmonic RF system for bunch compression.

Coupled bunch instabilities arising in the case of loss of Landau Damping may also be suppressed by a coupled bunch feedback system preserving the longitudinal emittance. In this case the bunches are shorter than in the case where they blow-up due to instabilities. Estimates for the minimum required RF kick strengths for such a feedback system are given.

INTRODUCTION

In the present operation scenario for the LHC the longitudinal emittance is increased in a controlled way during acceleration by a factor of about 2.5 [2, 3], resulting in a full (4 $\sigma$) bunch length of 0.31 m at top energy. Because of this blow up, there will always be sufficient Landau damping to stabilize the beam, in particular longitudinal coupled bunch instabilities are suppressed.

To obtain shorter bunches at top energy one may switch off the controlled emittance blow up (scenario i). By doing so, the beam may no longer be sufficiently stabilized and we have to consider additional measures for the suppression of longitudinal coupled bunch instabilities.

Conceivable measures for the beam stabilization are the increase of the bunch to bunch frequency spread by RF amplitude modulation (scenario ii) and a coupled bunch feedback system.

The additional voltage of a higher harmonic RF system leads to a bunch compression and supplies automatically more Landau damping due to the shorter buckets (scenario iii).

BUNCH LENGTH WITH COMPLETELY PRESERVED EMITTANCE

The longitudinal emittance $\epsilon$ depends approximately on the bunch length $l$ as [4]

$$\epsilon = \pi \Delta t \Delta E$$

with [5]

$$\Delta t = \frac{l}{2c}$$

$$\Delta E = \beta \sqrt{\frac{2E_0}{\eta}} \frac{eV}{2\pi} \frac{1}{h} \left( 1 - \cos \left( \frac{\omega_{rt} l}{2e} \right) \right)$$

where $c$ is the speed of light, $\Delta E$ the energy deviation, $\Delta t$ the time deviation of a particle with phase deviation $\Delta \phi = \omega_{rt} l/2$, $\beta = \frac{c}{V}$, $E_0$ the energy of the synchronous particle and $\eta = \frac{1}{\gamma^2} - \frac{1}{\gamma}$ the slip factor. The slip factor follows from $\gamma$-transition $\gamma_1$, the factor $\gamma = \frac{E_r}{E_0}$ and the proton rest energy $E_0 = m_p c^2$.

In the case of LHC the slip factor is at top energy $\eta_{7 \text{TeV}} = 0.0003225$. Due to the low $\gamma$-transition of 55.678 $\text{GeV}$ it has approximately the same value at injection energy ($\eta_{450 \text{GeV}} = 0.003182$). The RF parameters of LHC are a sum voltage of $V = 16 \text{ MV}$ and the frequency of $\omega_{rt} = 2 \pi 400.8 \text{ MHz}$ resulting in the harmonic number $h = 35640$.

At the injection energy of $E_0 = 450 \text{ GeV}$ the longitudinal (4 $\sigma$) emittance will be $\epsilon = 1 \text{ eVs}$. It is increased during acceleration to $\epsilon = 2.5 \text{ eVs}$ at top energy $E_0 = 7 \text{ TeV}$ corresponding to a full bunch length of 0.31 m. With a completely preserved longitudinal emittance of $\epsilon = 1 \text{ eVs}$ the bunch length at top energy would be

$$l_{1 \text{ eVs,7 TeV}} = 0.19 \text{ m}.$$  

MINIMUM BUNCH LENGTH STABILIZED BY LANDAU DAMPING

According to [6] the coherent frequency shift $\Delta f_{\text{imp}}$ caused by the effective impedance $Z_L/n$ is at the LHC given by

$$\Delta f_{\text{imp}} = 0.0237 \frac{\text{Hz}}{\Omega} \left( \frac{1}{n} \frac{Z_L}{n} \right) \left( \frac{N_0}{10^{21}} \right) \left( \frac{l}{\text{m}} \right)^{-3}.$$

67
\(N_b\) is the number of particles per bunch and \(l\) the bunch length. This relation is valid for the bunch spacing of 25 ns.

Landau damping will suppress instabilities up to a threshold value of

\[
\Delta f_{\text{threshold,Landau}} = 6.45 \text{ Hz} \left( \frac{l}{m} \right)^2 .
\]

By setting \(\Delta f_{\text{threshold,Landau}} = \Delta f_{\text{imp}}\) we can calculate the minimum bunch length at which Landau damping will be lost:

\[
\frac{l_{\text{min}}}{m} = \sqrt{3.57 \cdot 10^{-3} \frac{\Omega}{\text{Im} \frac{Z_L}{n}}} \left( \frac{\text{Im} \frac{Z_L}{n}}{N_b} \right)^{10^{11}}
\]

Table 1 shows the minimum bunch length for the most recent value of the longitudinal impedance of \(\text{Im} \frac{Z_L}{n} = 0.08 \Omega\) [3] and an older estimate \(\text{Im} \frac{Z_L}{n} = 0.28 \Omega\) [6] which may be viewed as a worst case value.

Table 1: Minimum bunch length at which Landau damping will be lost without additional measures.

<table>
<thead>
<tr>
<th>(N_b)</th>
<th>(\text{Im} \frac{Z_L}{n})</th>
<th>(0.08 \Omega)</th>
<th>(0.28 \Omega)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.1 (\cdot 10^{11})</td>
<td>0.20 m</td>
<td>0.26 m</td>
<td></td>
</tr>
<tr>
<td>1.7 (\cdot 10^{11})</td>
<td>0.22 m</td>
<td>0.28 m</td>
<td></td>
</tr>
</tbody>
</table>

For ultimate intensity \(N_b = 1.7 \cdot 10^{11}\) and also for nominal intensity \(N_b = 1.1 \cdot 10^{11}\) Landau damping will be lost by switching off the controlled emittance blow up (scenario i) because the minimum bunch length determined by the impedance exceeds the bunch length of \(l_{1 \text{eV}_\text{eVs},7 \text{TeV}} = 0.19 \text{ m}\) following from the initial longitudinal emittance.

### STABILIZATION BY RF AMPLITUDE MODULATION

The minimum bunch length values for \(\frac{Z_L}{n} = 0.08 \Omega\) in Table 1 are not too far away from the bunch length value for an emittance of \(e = 1 \text{ eVs}\). Hence, a method to double the instability threshold \(\Delta f_{\text{threshold}}\) may be sufficient to stabilize the beam. Such a method is the modulation of the RF amplitude (scenario ii). For the LHC it increases the instability threshold by about [7]

\[
\Delta f_{\text{threshold,AM}} = 0.51 \text{ Hz}
\]

resulting in

\[
\Delta f_{\text{threshold}} = \Delta f_{\text{threshold,Landau}} + \Delta f_{\text{threshold,AM}} .
\]

For \(\Delta f_{\text{threshold}} = \Delta f_{\text{imp}}\) the minimum bunch length at which Landau damping will be lost is given in Table 2.

The modulation of RF amplitude should be sufficient for stabilizing the beam for \(\text{Im} \frac{Z_L}{n} = 0.08 \Omega\) and a shorter bunch length is possible. Dividing the minimum bunch length for a given emittance by the minimum bunch length for a given impedance results in the safety margins of 26% and 12% respectively. Those safety margins are smaller than the safety margins of 55% and 41% obtained at normal operation with artificial beam blow up.

Table 2: Minimum bunch length at which Landau damping will be lost with rf amplitude modulation.

<table>
<thead>
<tr>
<th>(N_b)</th>
<th>(\text{Im} \frac{Z_L}{n})</th>
<th>(0.08 \Omega)</th>
<th>(0.28 \Omega)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1.1 (\cdot 10^{11})</td>
<td>0.15 m</td>
<td>0.21 m</td>
<td></td>
</tr>
<tr>
<td>1.7 (\cdot 10^{11})</td>
<td>0.17 m</td>
<td>0.23 m</td>
<td></td>
</tr>
</tbody>
</table>

In case the initial emittance is smaller or the longitudinal impedance larger, additional measures have to be taken.

### HIGHER HARMONIC RF SYSTEM

A RF system with three times the fundamental RF frequency (3 \(\times\) 400 MHz = 1.2 GHz) [1] for bunch compression is under investigation [9] (scenario iii). With the resulting double RF system one injects into the buckets of the 400 MHz system, whereas the 1.2 GHz system is operated such, that its contribution to the sum voltage is 1/4 or less. During acceleration the voltage of the 1.2 GHz system is increased such, that it completely builds up the buckets at top energy.

The energy deviation \(\Delta E\) is in the non-accelerating case given by

\[
\Delta E = \beta \sqrt{\frac{2 E_s}{\eta} \left( \frac{e}{2\pi} \left( \frac{V_{100}}{h_{100}} \left( 1 - \cos \left( \frac{\omega_{12} t}{2e} \right) \right) \right)^3 + \frac{V_{1200}}{3 h_{1200}} \left( 1 - \cos \left( \frac{3 \omega_{12} t}{2e} \right) \right)^3 \right)^{-1/2}}
\]

resulting with the voltages \(V_{s00} = 16 \text{ MV}\) and \(V_{1200} = 43 \text{ MV}\) in the bunch length of

\[
l_{1 \text{eV}_\text{eVs},7 \text{TeV}} = 0.12 \text{ m},
\]

in the case the emittance is completely preserved. In case some (controlled) blow up will take place one may obtain the in [1] proposed bunch length value of

\[
l_{1.8 \text{eV}_\text{eVs},7 \text{TeV}} = 0.16 \text{ m}.
\]

By operating the LHC at top energy with both RF systems the effective voltage in the bucket core becomes \(V_{200,\text{eff}} = 48.3 \text{ MV}\) resulting in a synchrotron frequency of \(f_s = 72 \text{ Hz}\) (\(f_s = 23.9 \text{ Hz}\) for single RF). These values are obtained from \(f_{0} \propto \sqrt{h V}\) and \(f_s = \sqrt{1 + \frac{r_{12}^2}{h_{12}^2}} f_{0}\), where \(r_{12} = V_{400}/V_{1200}\) is the voltage ratio and \(h_{12} = 1/3\) the ratio of the harmonic numbers [10, p. 76].
With $\Delta f_{imp} \propto f_b / (h V)$ [6] the coherent frequency shift obeys

$$\Delta f_{imp} = 0.0079 \frac{Hz}{\Omega} \left( \frac{\text{Im} \, Z_1}{m} \right) \left( \frac{N_0}{10^{11}} \right) \left( \frac{l}{m} \right)^{-3}.$$  

Considering $\Delta f_{\text{threshold,Landau}} \propto f_b$ an approximation for the Landau damping in a double RF system is for $r_{12} < 1 / h_{12}$ given by [10, p. 77]

$$\Delta f_{\text{threshold,Landau}} = 19.4 \, \text{Hz} \, \frac{1 + r_{12} h_{12}^2}{\sqrt{1 + r_{12} h_{12}}} \left( \frac{3L}{m} \right)^2.$$  

The fact that the higher harmonic RF system mainly defines the bucket area at top energy is taken into consideration by the choice of $r_{12} = V_{1400} / V_{1200}$ and the factor three in front of the bunch length.

At top energy the condition $r_{12} < 1 / h_{12}$ is fulfilled and we get

$$\Delta f_{\text{threshold,Landau}} = 167 \, \text{Hz} \left( \frac{l}{m} \right)^2.$$  

Table 3: Minimum bunch length at which Landau damping will be lost for a double rf system.

<table>
<thead>
<tr>
<th>$l_{\text{min}}$</th>
<th>$\text{Im} , \frac{Z_1}{m} = 0.08 , \Omega$</th>
<th>$0.28 , \Omega$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$N_b = 1.1 \cdot 10^{11}$</td>
<td>0.08 m</td>
<td>0.11 m</td>
</tr>
<tr>
<td>$N_b = 1.7 \cdot 10^{11}$</td>
<td>0.09 m</td>
<td>0.12 m</td>
</tr>
</tbody>
</table>

Table 3 shows the minimum bunch length at which Landau damping will be lost at the double RF system. Only in the case of an unexpected large effective impedance, Landau damping may be lost for beams with ultimate intensity and completely preserved emittance.

The estimates presented in Table 3 contain some uncertainties: The higher harmonic RF system itself will contribute to the effective longitudinal impedance, its contribution is under investigation [9]. Furthermore, shorter bunches lead in general to slightly larger effective impedances, depending on the exact frequency distribution of the accelerator impedance. In the estimates presented here, this fact is not taken into account.

Even in the case of an unexpected high effective impedance the beam should be sufficiently stabilized by a controlled emittance blow up to 1.8 eVs by Landau damping takes place. Hence, the required kick voltage $V_{\text{THK}}$ of a longitudinal coupled bunch feedback depends on the maximum growth rate $\Delta f_{\text{threshold,Landau}} / f_b$ stabilized by Landau damping, the RF voltage $V_{\text{RF}}$ and the minimum detectable phase oscillation $\Delta \phi_{\text{det}}$, like [6]

$$V_{\text{THK}} > 2 \, V_{\text{RF}} \frac{\Delta f_{\text{threshold,Landau}}}{f_b} \frac{\Delta \phi_{\text{det}}}{\phi}.$$  

For the single harmonic RF system the synchrotron frequency is at top energy $f_b = 23.9 \, \text{Hz}$. A reasonable choice for $\Delta f_{\text{threshold,Landau}}$ is the value corresponding to the emittance where the beam is reliably stabilized by Landau damping only.

If the bunch phase oscillations are detected in the same way as in the HERA proton ring [10, p. 32] the minimum detectable phase oscillation is $\Delta \phi_{\text{det}} \approx 2 \pi / 500 = 0.2^\circ$. Digital filters may decrease this value. Using $\Delta f_{\text{threshold,Landau}}$ for an emittance of $\epsilon = 2.5$ eVs results in the minimum required kick voltage for a single RF system of

$$V_{\text{THK}} > 3 \, \text{kV}.$$  

In the case of the double RF system, the beam is expected to be stabilized by Landau damping only, for longitudinal emittances of 1 eVs.

This estimate does not take into account transient effects at injection or during acceleration. Fast damping of synchrotron oscillations caused by such effects may require a larger kick strength.

**SUMMARY**

With the given RF system, bunches shorter than about 25 cm are not possible in the LHC by switching off the controlled emittance blow up used in standard operation (scenario i). For the most recent estimate of the effective longitudinal impedance $\text{Im} \, \frac{Z_1}{m} = 0.08 \, \Omega$, a modulation of the RF phase may be sufficient to stabilize the beam for the case of controlled emittance blow up (scenario ii). This is no longer the case for initial emittances smaller than 1 eVs or an impedance larger than 0.08 $\Omega$. An impedance larger than 0.08 $\Omega$ may be caused by collimators for machine protection and for the reduction of background events in the high energy experiments ob by other operational constraints. In such a case additional measures have to be taken.

One measure is a higher harmonic RF system for bunch compression (scenario iii). For $\text{Im} \, \frac{Z_1}{m} = 0.08 \, \Omega$ the beam should be well stabilized by Landau damping. Investigations on the contribution of this RF system itself to the effective impedance are under way [9].

Alternatively one may set up a longitudinal coupled bunch feedback system. Then, the minimum bunch length obtained by suppressing coupled bunch oscillations depends on the initial emittance and the (noise) performance of the feedback system.
Table 4 gives an overview on the minimum bunch length in LHC for a single harmonic RF system, a double harmonic RF system and the measures required for suppressing longitudinal coupled bunch instabilities.

Table 4: Minimum bunch length in an upgraded LHC.

<table>
<thead>
<tr>
<th>RF system</th>
<th>‘low’</th>
<th>‘high’</th>
<th>( l_{\text{BR}} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>400 MHz</td>
<td>RF AM</td>
<td>feedback</td>
<td>0.19 m</td>
</tr>
<tr>
<td>400 MHz &amp; 1.2 GHz</td>
<td>blow up</td>
<td>blow up</td>
<td>0.16 m</td>
</tr>
<tr>
<td></td>
<td></td>
<td>RF AM</td>
<td>0.12 m?</td>
</tr>
</tbody>
</table>

A set up of fast multi bunch beam diagnostics and data logging, comparable to the system used at the HERA proton ring [8], may be required for the study of the instabilities and the supervision of the measures taken.

Longitudinal microwave instabilities are expected to appear for bunches shorter than about 5.5 cm [1]. This is less than 1/2 of the minimum bunch length discussed in this paper, resulting in a minimum safety margin of more than 100%. For that reason, longitudinal microwave instabilities have not been treated in this paper.

ACKNOWLEDGMENTS

I would like to thank Thomas Bohl for having a look at the contents of this workshop contribution and suggestions.

REFERENCES


MORPHOLOGICAL AND STRUCTURAL STUDIES OF CRYSTALS FOR CHANNELING OF RELATIVISTIC PARTICLES

S. Baricordi, V. Guidi, C. Malagù, G. Martinelli, E. Milan, M. Stefanich, Department of Physics and INFN, Via Paradiso 12, I-44100 Ferrara, Italy
A. Carnera, A. Sambo, C. Scian, A. Vomiero\(^1\), Department of Physics, Via Marzolo 8, I-35131 Padova, Italy
G. Della Mea, S. Restello, INFN Laboratori Nazionali di Legnaro, Viale Università 2, I-35020 Legnaro (PD), Italy
V.M. Biryukov, Yu.A. Chesnokov, Institute for High Energy Physics, Protvino, Russia
Yu. M. Ivanov, Petersburg Nuclear Physics Institute, Gatchina, 188350, Russia
W. Scandale, CERN, Geneva 23, CH-1211, Switzerland

Abstract

Channeling of relativistic particles through a crystal may be useful for many applications in accelerators. It has been experimented that significant role is played by the method used to clean the crystal surfaces that the particles beam encounters first. It was experimentally observed that chemical cleaning of the surfaces via etching leads to better performance than conventional mechanically treated samples do. We investigated the physical reasons for such a behaviour through characterization of the surfaces of the crystal. We observed that mechanical dicing causes a superficial layer rich in dislocations and lattice imperfections, extending tens of microns into the crystal, which is removed via etching. Such a disordered layer is the first portion of the crystal experienced by the incoming particles and results in a mosaicity degree that exceeds the critical angle of relativistic particles for channeling, i.e., it acts like an amorphous layer.

INTRODUCTION

Channeling of relativistic charged particles would be a powerful method in accelerator physics [1-3] because it may be an alternative for construction of the components presently used in accelerators such as dipoles, collimators and focusing elements [4]. For this reason, the use of bent crystals for beam extraction in circular accelerators has been intensively investigated in several laboratories during last decade. The advantages of the method of channeling are low cost, compactness and minimal disturbance to the beam in the environment of the accelerator.

Silicon is an excellent candidate to fabricate crystals for channeling because of its high crystalline perfection, low cost and established knowledge of its handling. Since pioneering experiments [5], a continuous improvement in performance has been carried out, arising to new schemes for the crystals and also to technological development.

\(^1\) Also INFN-LNL

Building Si crystals with the right dimensions for each specific application does demand dicing, lapping, and other operations that alter the original quality of the lattice. Preparation of samples induces a “dead layer” at surface, characterized by a great number of defects and crystalline disorder, which does not act as an active layer for channeling. It has been demonstrated that significant improvement in the features of the samples for channeling can be achieved once a specific etching treatment is imparted to the crystal surface [6].

In this work we report on a systematic analysis of the crystalline perfection of the surface of Si crystals applied in channeling experiments, depending on the preparation methodology.

Figure 1: Extraction efficiency for 70-GeV protons. Recent results [6] (*, strips, 1.8, 2.0, and 4 mm long), results of 1999-2000; (○, “O-shaped” crystals 3 and 5 mm), and of 1997 (⊗, strip 7 mm). Also shown (○) is the Monte Carlo prediction for a perfect crystal with 0.9 mrad bending.

CHANNELING EFFICIENCY

In the frame of the CERN-INTAS collaboration 2000-132, we recently achieved a substantial progress with Si crystal-assisted beam deflection at the 70-GeV
accelerator at IHEP. Extraction efficiency of the order of 85% was repeatedly obtained for an impinging intensity as high as \(10^{12}\) protons [6].

Key reason of this successful operation was the use of very short crystals for extraction. Crystal length was selected close to the optimal value foreseen by the physics of proton channeling. Fig. 1 shows theoretical predictions of extraction efficiency as a function of the crystal length, for an impinging proton beam of 70-GeV, as compared to experimentally recorded levels with crystals of different size and design.

A comparative study [7] about the performance of mechanically polished samples vs. chemically etched specimens was carried out at IHEP. Protons, 70-GeV in energy, were extracted from the main stream through both types of crystals. Fig. 3 shows the results of the tests for the chemically polished deflectors vs. those for the unpolished samples. The crystals with chemically polished faces have shown the best perform for beam extraction. The profile of the beam bent by a chemically etched crystal turned out to be more uniform and sharp. Its width was ascribed to the crystal thickness once beam divergence within the critical angle of channeling (20 \(\mu\)rad at that energy) was taken into account. On the other hand, the beam bent by the mechanically polished crystal exhibited irregularities, corresponding to an angular spread of the order of 100 \(\mu\)rad.

**MORPHOLOGICAL CHARACTERIZATION**

As a general rule for microelectronics, when the dimensions of silicon devices scale down, wafers with minimal impurities are needed, thus great care must be taken for cleaning procedures. Bearing this scheme in mind, we borrowed from microelectronics a cleaning methodology to manufacture the samples for channelling.

Figure 2: Scheme of the bent crystal plate. Dimensions are expressed in mm, \(y\) is the direction tangent to the incident beam on the middle of the crystal, \(x\) is oriented along the thickness (crystalline direction (111) of silicon) and \(z\) is oriented along the height. The crystals were manufactured at the Semiconductors and Sensors Laboratory of the University of Ferrara as narrow strips, about 2-mm thick in the direction of the beam.

A recent design to produce a short crystals is the use of the anisotropic properties of a crystal lattice. From the theory of elasticity it is known that bending a crystal plate in the longitudinal direction causes "anticlastic bending" or twists to appear in the orthogonal direction. For Si (111)-oriented, the crystal plate takes the shape of a saddle as sketched in Fig. 2.

Figure 3: Image of the beam deflected through \(2\) mechanically treated (left) and chemically polished crystals (right). The profile of the beam bent by a chemically polished crystal is more uniform and sharp (20 \(\mu\)rad vs 100 \(\mu\)rad at 70 GeV).

Figure 4: (a) A sample cut to form 0.497\(\times\)20\(\times\)5 mm\(^3\) (thickness, length, height respectively) by means of a mechanical dicing saw with a 30000 rpm angular velocity and 0.150 mm thickness of diamond blade. The linear velocity was 2 mm/sec. (b) A mechanically treated sample and (c) a chemically polished sample.

The first step is removal of organic and metallic impurities that are present on the wafer surface. The process of dicing a wafer by a diamond-blade saw to reach the wanted dimensions induces a superficial layer rich in scratches, dislocation, line defects and anomalies. In order to remove the defects, two types of surface cleaning are carried out: mechanical polishing and chemical etching. For mechanical polishing, the sample is fixed on a special slide, which is put onto a rotating plate covered with different abrasive cloths. The chemical method consists of a planar etching [7], which allows the removal of planes of silicon one by one. Fig. 4 consists of electron-microscopy photos of some samples: just after the cut (a), mechanically
treated (b) and chemically polished (c). Fig. 4(a) shows a sample cut by the dicing saw at 30000 rpm angular velocity by a 150-µm-thick diamond blade. The linear velocity was 2 mm/s. Fig. 4(b) illustrates a mechanically-polished sample after a sequence of four abrasive cloths with decreasing roughness. Fig. 4(c) is a chemically polished sample through 30 min chemical-etching. Chemical etching allows slow erosion for precise manufacturing of the samples to an extent that is not otherwise possible.

Atomic force microscopy (AFM) was carried out first to measure the standard roughness, \( R_a \), of the surfaces induced by dicing at top speed. An area of 40×40 µm\(^2\) was chosen for collection of images (Fig. 5) prior to (a) and after (b) chemical etching. Fig. 5(a) shows the levels of roughness as a function of the state of the surface.

Mechanical polishing highly reduces the roughness by a factor of five with respect to the unpolished surface. Rather surprisingly, 20 min of chemical etching enhances the roughness even if compared to the as-cut sample. However, for longer etching times (up to 40 min), \( R_a \) tends to decrease (see Fig. 6(b)).

![AFM images of the surface of an as-diced Si crystal](image1)

Figure 5: AFM images of the surface of an as-diced Si crystal (a) and after 40 min chemical etching (b).

![RBS-channeling spectra](image2)

Figure 7: RBS-channeling spectra recorded using alpha particles (a) and protons (b). Reference spectra from a “perfect” crystal in random and best channeling orientation are reported. For both \( \alpha \) particles and protons the yield of the samples after chemical etching (E) is the same as that of the reference crystal under channeling regime. The as-cut sample diced at higher speed (5 C), instead, exhibits higher yield with respect to the sample cut at 0.5 mm/min (0.5 C).

In Fig. 7(a), RBS-channeling spectra obtained from 2.0-MeV \(^3\)He\(^+\) beam on sample cut at 0.5 mm/min are reported after normalization to the collected charge. Two reference spectra of a random and aligned Si crystal are also reported. Under such experimental conditions, we estimated an in-depth analyzed region of about 1.5 µm. Three main observations can be drawn. First: a surface Si peak is recorded in all the spectra due to surface scattering; the peaks for the reference crystal and for sample E (chemical etching) are similar, while samples C (as-cut) and M (mechanically polished, not reported) exhibit a longer tail towards the bulk of the crystal. This evidence highlights the presence of a
surface disordered structure for samples M and C, which is not detected for sample E. Moreover, the extension of the unchanneled region for sample M is broader than for sample C. Second: irrespective of the surface peak, the RBS signal from the inner part of the sample exhibits a very regular shape, indicating a homogeneous in-depth distribution of defects, which are responsible for dechannealing events as far as the analyzed depth is concerned. Third: the normalized height of spectra pertaining to samples M and C is the same at any depth but the surface region, and is systematically five times higher than the height for sample E. This evidence clearly shows that, despite the enhanced roughness induced by 30 min of chemical etching, the etched crystal exhibits highly improved crystalline order with respect to both C and M samples.

A complete analysis on the entire series of samples was carried out using a 2.0 MeV proton beam, which is more penetrating and allows investigation of a deeper region (up to 14 μm) with respect to α particles (see Fig. 7(b)). The as-diced crystal cut at 5 mm/min exhibits a higher yield for backscattered protons, under channeling conditions, with respect to the as-diced crystal cut at 0.5 mm/min (Fig. 7(b)), i.e., the samples cut at higher speed highlighted larger crystalline disorder. Mechanically polished samples show an intermediate yield with respect to sample C and E, though they exhibited little reproducibility. Concerning chemical etching, backscattering yield of all samples was the same irrespective of the dicing speed and perfectly overlaps with the spectrum of a reference single-crystal. This is a clear proof of the effectiveness of chemical etching for removal of the superficial damaged layer, at least within the penetration range of the technique.

CONCLUSIONS

Considerable progress in crystal channeling research has been obtained over the years. Extraction efficiency of the order of 85% was repeatedly obtained for an impinging intensity as high as 10^{12} protons. Channeling in crystals proved to be positively affected by superficial etching treatments. Chemically polished samples exhibit a more ordered crystalline structure, indicating the removal of the amorphized layer, which was formed when the crystal was diced.

REFERENCES

Limitations due to Electron-Cloud Heat Load for the LHC and its Upgrade

D. Schulte, F. Zimmermann, CERN, Geneva, Switzerland

Abstract

In the LHC beam pipe, an electron cloud can build up due to photoemission from synchrotron radiation and via a beam-induced multipacting process. This electron cloud gives rise to an additional heat load on the beam screen which protects the cold bore of the superconducting magnets. The heat load from the electron cloud depends on beam parameters such as the bunch intensity, the number of bunches, the bunch length, and the bunch spacing, as well as on surface parameters, in particular the maximum secondary emission yield and the elastic reflectivity of low-energy electrons. The maximum tolerable heat load is determined by the available cooling capacity. Simulations of the electron-cloud induced heat load in various LHC configurations yield the limiting values of secondary emission yield and reflectivity which are required at 25-ns bunch spacing to reach the nominal and ultimate LHC parameters, respectively, at 12.5-ns bunch spacing for a ‘conventional upgrade’, and at 75-ns bunch spacing for an upgrade in the regime of a large Piwinski parameter.

1 SCENARIOS

We consider the four scenarios listed in Table 1: the nominal LHC, the ultimate, the baseline upgrade with more and shorter bunches, and the alternative upgrade in a regime of a large Piwinski parameter [1]. The longitudinal bunch profile is Gaussian except for the last case, where a uniform profile is assumed.

Other parameters adopted in the electron-cloud simulations are summarized in Table 2. The photoemission rate corresponds to that expected after surface conditioning [2, 3]. In addition, it is assumed that 20% of the primary photons are reflected, with a distribution of the form \( \cos^2 \phi \) [2, 4], where \( \phi \) describes the angle of reflected photons with respect to the horizontal plane as viewed from the horizontally outward primary impact point of the synchrotron radiation.

2 SIMULATION RESULTS

Figure 1 presents simulated average arc-cell heat loads in the four scenarios as a function of the maximum secondary emission yield \( \delta_{\text{max}} \) for three different values of the low-energy electron reflectivity \( R \). The averaging over an arc cell was performed using as weights the relative lengths of drifts, quadrupoles, and dipole regions; the values for the sextupoles were taken to be equal to those of the quadrupoles.

The simulated heat loads should be compared with the available cooling capacity, which is of order 1–2 W/m [5]. From Fig. 1, we deduce that for the nominal and ultimate LHC, a value of \( \delta_{\text{max}} \) smaller than about 1.5 must be reached, which has been routinely achieved in the CERN SPS [6]. For the baseline upgrade, \( \delta_{\text{max}} \) needs to decrease below 1.2, which has also been demonstrated in the laboratory, e.g., [7], but not yet in an accelerator. The Piwinski upgrade does not pose any challenge from the electron-cloud viewpoint, since all values of \( \delta_{\text{max}} \) up to 2.0 would be acceptable.

Figure 2 shows the same results as Fig. 1 converted into heat-load contours in the \( \delta_{\text{max}}-R \) plane.

Simulated average electron densities, at the center of the vacuum chamber, are displayed in Fig. 3. The typical threshold of the electron-driven fast transverse-mode coupling instability in the LHC (albeit at injection, whereas the electron-cloud build-up simulations are — pessimistically — performed at top energy) is found for densities of about \( 5 \times 10^{11} \text{ m}^{-3} \) [10]. At the ultimate LHC the simulated central electron density is lower than for the nominal LHC. This is consistent with the theoretical expectation [8, 9, 1]. The \( \delta_{\text{max}} \) values required to guarantee beam stability at injection are less than 1.4 for the nominal, less than 1.5 or 1.6 for the ultimate, less than 1.5 for the baseline upgrade, and not restricted for the Piwinski upgrade.

Figure 4 shows the simulated flux of electrons incident on the wall with energies larger than 30 eV. These electrons contribute to gas desorption and to surface conditioning. Typical fluxes near the limiting values of \( \delta_{\text{max}} \) set by the heat load are between 1 and 10 mA per meter for the nominal, ultimate and baseline upgrade. In the case of the Piwinski upgrade, the flux is only a few 100 \( \mu \text{A} \) per meter.

3 CONCLUSIONS

A maximum secondary electron emission yield below 1.5, as routinely achieved in the SPS, will support the nominal LHC intensity and also the ultimate. The surface requirements for the ultimate LHC are even slightly relaxed compared with the nominal case. The upgrade towards shorter bunch spacing demands a reduced secondary emission yield of less than 1.2. For the alternative upgrade path with fewer and longer bunches, no limitation from electron cloud is expected. The requirements imposed by the heat load are comparable to, or slightly more stringent than, those from the beam stability via the central electron density.
Table 1: Beam parameters

<table>
<thead>
<tr>
<th>scenario</th>
<th>nominal</th>
<th>ultimate</th>
<th>baseline upgrade</th>
<th>‘Piwinski upgrade’</th>
</tr>
</thead>
<tbody>
<tr>
<td>no. bunches</td>
<td>2808</td>
<td>2808</td>
<td>5616</td>
<td>936</td>
</tr>
<tr>
<td>bunch population</td>
<td>11.5 × 10^{10}</td>
<td>17 × 10^{10}</td>
<td>17 × 10^{10}</td>
<td>60 × 10^{10}</td>
</tr>
<tr>
<td>bunch spacing</td>
<td>25 ns</td>
<td>25 ns</td>
<td>12.5 ns</td>
<td>75 ns</td>
</tr>
<tr>
<td>bunch length</td>
<td>7.55 cm (σ)</td>
<td>7.55 cm (σ)</td>
<td>3.78 cm (σ)</td>
<td>50 cm (full)</td>
</tr>
<tr>
<td>bunch profile</td>
<td>Gaussian</td>
<td>Gaussian</td>
<td>Gaussian</td>
<td>uniform</td>
</tr>
<tr>
<td>rms transv. beam size σ_{x,y}</td>
<td>300 μm</td>
<td>300 μm</td>
<td>300 μm</td>
<td>300 μm</td>
</tr>
<tr>
<td>chamber half aperture h_{x,y}</td>
<td>22, 18 mm</td>
<td>22, 18 mm</td>
<td>22, 18 mm</td>
<td>22, 18 mm</td>
</tr>
</tbody>
</table>

Table 2: Electron-cloud simulation parameters

<table>
<thead>
<tr>
<th>variable</th>
<th>symbol</th>
<th>value</th>
</tr>
</thead>
<tbody>
<tr>
<td>photo-emission rate / e^+ / meter</td>
<td>d\lambda_{pe} / ds</td>
<td>5.08 × 10^{-4}</td>
</tr>
<tr>
<td>photon reflectivity</td>
<td>R</td>
<td>20%</td>
</tr>
<tr>
<td>reflected photon distr.</td>
<td>dN_{refl,-\gamma} / d\phi</td>
<td>\cos^{2} \phi</td>
</tr>
<tr>
<td>max. secondary emission yield</td>
<td>\delta_{max}</td>
<td>variable</td>
</tr>
<tr>
<td>energy at max. sec. em. yield</td>
<td>\epsilon_{max}</td>
<td>varying with \delta_{max}</td>
</tr>
</tbody>
</table>

Figure 1: Simulated average arc heat load as a function of maximum secondary emission yield \( \delta_{max} \) for the nominal LHC (top left), the ultimate (top right), the baseline upgrade (bottom left), and the upgrade with large Piwinski angle (bottom right). The three curves refer to three different values for the reflectivity of low-energy electrons (0, 50% and 100%).
Figure 2: Contour lines of constant heat load in W/m as a function of maximum secondary emission yield (horizontal axis) and reflectivity of low energy electrons (vertical axis), for the nominal LHC (top left), the ultimate (top right), the baseline upgrade (bottom left), and the upgrade with large Piwinski angle (bottom right).
Figure 3: Simulated central electron density in m$^{-3}$ as a function of maximum secondary emission yield $\delta_{\text{max}}$, for the nominal LHC (top left), the ultimate (top right), the baseline upgrade (bottom left), and the upgrade with large Piwinski angle (bottom right). The three curves refer to three different values for the reflectivity of low-energy electrons (0, 50% and 100%).
Figure 4: Simulated electron flux at the wall (with electron energy larger than 30 eV) in m$^{-1}$A as a function of maximum secondary emission yield $\delta_{\text{max}}$, for the nominal LHC (top left), the ultimate (top right), the baseline upgrade (bottom left), and the upgrade with large Piwinski angle (bottom right). The three curves refer to three different values for the reflectivity of low-energy electrons (0, 50% and 100%).
4 ACKNOWLEDGMENTS

We thank V. Baglin and F. Ruggiero for helpful discussions and motivating some of these studies.

5 REFERENCES


A LONG STRAIGHT SECTIONS

We here report electron-cloud build-up simulations for the long straight sections of the nominal LHC. Considering a single beam in a chamber with 80 mm diameter, simulations are run for $\delta_{\text{max}} = 1.1$ and $\delta_{\text{max}} = 1.4$ as well as for three different values of low-energy electron reflectivity ($R = 0, 0.5$ and $1.0$). Computed are the heat load, the electron flux and the average electron energy. The latter two quantities are important for estimating the dynamic vacuum pressure and the scrubbing efficiency.

We assume the standard synchrotron radiation flux and photo-emission yields as expected for the arcs after scrubbing. This is very pessimistic, since the estimated average photon flux in the straight sections is about 20 times smaller than that in the arcs [11]. The beam is represented as a round beam centred in a round chamber, with an rms beam size of 0.3 mm in both transverse planes, which is the typical beam size in the arcs. We have verified that for a 5 times larger beam size the results change by a few percent at the most. No external magnetic field was included.

Results are summarized in Tables 3 and 4, for the nominal bunch intensity and a 33% The typical electron flux is of the order $10^{17}$ electrons per meter per second, the average electron energy varies between 100 and 200 eV, and the heat load can be up to a few W/m. 

80
Table 3: Results for the long straight section at the nominal intensity of $1.15 \times 10^{11}$ protons per bunch. The flux is quoted per unit length; to obtain the flux per unit beam-pipe area in ($\text{m}^{-2}$) the numbers must be multiplied by about a factor 4.

<table>
<thead>
<tr>
<th>sec. emission parameters</th>
<th>power</th>
<th>total flux</th>
<th>flux ($E &gt; 30$ eV)</th>
<th>mean $e^-$ energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\delta_{\text{max}} = 1.4, R = 0.0$</td>
<td>3.72 W/m</td>
<td>$1.46 \times 10^{17}$ e$^-$/m/s</td>
<td>$1.20 \times 10^{17}$ e$^-$/m/s</td>
<td>159 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.4, R = 0.5$</td>
<td>4.50 W/m</td>
<td>$2.15 \times 10^{17}$ e$^-$/m/s</td>
<td>$1.64 \times 10^{17}$ e$^-$/m/s</td>
<td>131 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.4, R = 1.0$</td>
<td>6.55 W/m</td>
<td>$4.21 \times 10^{17}$ e$^-$/m/s</td>
<td>$2.72 \times 10^{17}$ e$^-$/m/s</td>
<td>97 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.1, R = 0.0$</td>
<td>0.97 W/m</td>
<td>$2.84 \times 10^{16}$ e$^-$/m/s</td>
<td>$2.66 \times 10^{16}$ e$^-$/m/s</td>
<td>213 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.1, R = 0.5$</td>
<td>1.25 W/m</td>
<td>$3.85 \times 10^{16}$ e$^-$/m/s</td>
<td>$3.55 \times 10^{16}$ e$^-$/m/s</td>
<td>203 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.1, R = 1.0$</td>
<td>1.25 W/m</td>
<td>$5.83 \times 10^{16}$ e$^-$/m/s</td>
<td>$5.21 \times 10^{16}$ e$^-$/m/s</td>
<td>185 eV</td>
</tr>
</tbody>
</table>

Table 4: Results for the long straight section at a reduced intensity of $0.77 \times 10^{11}$ protons per bunch. The flux is quoted per unit length; to obtain the flux per unit beam-pipe area in ($\text{m}^{-2}$) the numbers must be multiplied by about a factor 4.

<table>
<thead>
<tr>
<th>sec. emission parameters</th>
<th>power</th>
<th>total flux</th>
<th>flux ($E &gt; 30$ eV)</th>
<th>mean $e^-$ energy</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\delta_{\text{max}} = 1.4, R = 0.0$</td>
<td>0.327 W/m</td>
<td>$2.34 \times 10^{15}$ e$^-$/m/s</td>
<td>$2.19 \times 10^{15}$ e$^-$/m/s</td>
<td>87 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.4, R = 0.5$</td>
<td>0.633 W/m</td>
<td>$4.57 \times 10^{16}$ e$^-$/m/s</td>
<td>$4.12 \times 10^{16}$ e$^-$/m/s</td>
<td>87 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.4, R = 1.0$</td>
<td>1.219 W/m</td>
<td>$9.25 \times 10^{16}$ e$^-$/m/s</td>
<td>$7.49 \times 10^{16}$ e$^-$/m/s</td>
<td>82 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.1, R = 0.0$</td>
<td>0.093 W/m</td>
<td>$5.68 \times 10^{15}$ e$^-$/m/s</td>
<td>$5.20 \times 10^{15}$ e$^-$/m/s</td>
<td>103 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.1, R = 0.5$</td>
<td>0.126 W/m</td>
<td>$7.29 \times 10^{15}$ e$^-$/m/s</td>
<td>$6.72 \times 10^{15}$ e$^-$/m/s</td>
<td>108 eV</td>
</tr>
<tr>
<td>$\delta_{\text{max}} = 1.1, R = 1.0$</td>
<td>0.180 W/m</td>
<td>$1.01 \times 10^{16}$ e$^-$/m/s</td>
<td>$9.24 \times 10^{15}$ e$^-$/m/s</td>
<td>111 eV</td>
</tr>
</tbody>
</table>
Overview of Possible LHC IR Upgrade Layouts
J. Strait, N. Mokhov and T. Sen
Fermilab, PO Box 500, Batavia, IL 60510

Abstract
An upgrade of the LHC interaction regions could potentially increase the luminosity by a factor of two or more. Several IR layouts are presented. The challenges and open questions related to the optics design, energy deposition and magnet design are discussed.

I. LHC Luminosity Upgrade
The initial luminosity goal in the LHC is an ambitious $10^{34}$ cm$^{-2}$ sec$^{-1}$. The present plan is to achieve 10% of this luminosity in the first year of running, 30% in the second year and full luminosity in the fourth year [1]. After running at full luminosity for a few years, the need to reduce the statistical errors in the experimental data will require an upgrade in luminosity. This error is inversely proportional to the square root of the integrated luminosity and falls rapidly while the luminosity is increasing but falls very gradually once the luminosity plateaus.

Figure 1: Evolution of the luminosity, integrated luminosity, the Poisson error in arbitrary units and the time in years to halve this error.

Figure 1 illustrates the time evolution of the luminosity, the integrated luminosity, the Poisson error and the time to halve this error.

Assuming that the luminosity evolves as predicted, by the year 2012 it will take more than 7 years of running at the same luminosity to reduce the errors by a factor of two. By the middle of the next decade, the interaction region quadrupoles will be nearing the end of their expected radiation lifetime of 600-700 fb$^{-1}$, after having absorbed all the debris power from the collisions [2]. For these separate reasons, keeping the LHC physics program productive and the need to replace major components, we expect that an upgrade will be required in the years between 2012-2015. Of course the benefits of increasing the luminosity extend far beyond reducing the statistical errors. The increased physics potential includes extending the reach of electro-weak physics, the ability to study coupled vector gauge boson interactions, searching for new modes in super-symmetric theories (SUSY) and also searching for new massive objects, some of which could be manifestations of extra dimensions [3].

Increasing the luminosity by an order of magnitude to $10^{35}$ cm$^{-2}$ sec$^{-1}$ will be very challenging. No single path for success is guaranteed and we need to consider several ways to achieve this goal. Upgrading the IRs will almost certainly be one of the options taken. It is likely that a new class of superconducting magnets will be required and considerable R&D is necessary to demonstrate that these can be built. From past experience we know that this takes several years, and therefore the R&D is starting now.

II. IR Upgrades
Luminosity can be increased directly by lowering the $\beta^*$ at the IPs – this will require an upgrade of the IRs. Furthermore if the long-range beam-beam interactions are observed to severely impact beam lifetime during the first phase of the LHC, then the IR layout can be
changed to reduce the number of these interactions. An upgrade of the IR optics could feasibly reduce $\beta^*$ by 2-3 times, therefore a 10-fold increase in luminosity will also require an upgrade of beam parameters such as intensities and emittances. The IR layout must be equipped to deal with the resulting challenges of higher beam current and the 10-fold increase in power from the collision debris. The technical challenges associated with increased focusing in the IR are well known. The larger $\beta_{\text{max}}$ values in the IR magnets imply that excellent field quality and precise alignment of these magnets will be even more critical. However the principal challenge will arise from the energy deposited by the collision byproducts in these magnets. The power in these debris particles increases linearly with luminosity - an estimated 9kW per beam must be safely absorbed at a luminosity of $10^{35}$ cm$^{-2}$ sec$^{-1}$. This directly impacts the quench stability of the magnets due to the increase in local peak power density, increases the heat load on the cryogenic system and increases the radiation damage of components [2]. These challenges must be met by improved optics designs, detailed energy deposition calculations and new engineering designs for the magnets.

The baseline IR has quadrupoles placed as close as possible to the IP. This layout has the advantages of minimizing $\beta_{\text{max}}$ in the magnets for a given $\beta^*$, and modest peak power deposition in the magnets because quadrupoles are relatively inefficient at sweeping charged particles. The disadvantages are a significant number of parasitic interactions between the beams, and a common correction system within the IR that must act on both beams simultaneously.

Several layouts for an IR upgrade were proposed in Reference [1]. These designs assume that pole tip fields in the neighborhood of 13T, with 15-20% operating margin, will be achievable with the use of Nb$_3$Sn superconductor. Figure 2 shows the two most straightforward designs – one with quadrupoles first as in the baseline IR and the other with dipoles first. In both these designs, the crossing angle increases with $1/\sqrt{\beta^*}$ from the baseline optics. Magnets in both layouts start at 23m from the IP – as in the baseline. TAS absorbers are placed before the magnets in both designs. In the first layout, the quadrupoles have the same gradient but larger coil aperture, 110mm, rather than the 70mm aperture in the baseline.

![Figure 2: Straightforward IR upgrade layouts. Top (a): quadrupoles first, Bottom (b): dipoles first.](image-url)

This aperture is the largest feasible for the desired gradient. The magnet aperture sets an upper limit on the allowed $\beta_{\text{max}}$ – as remarked earlier, placing the quadrupoles closest to the IP leads to the smallest $\beta^*$ for this allowed $\beta_{\text{max}}$. The first separation dipole D1 in this dipole is made as short as possible to minimize the amount of power deposited at the end of the dipole furthest from the IP. In the second layout,
the beams are separated early with 10 m long dipoles D1 and D2. Space has been left for 5m
long neutral absorbers before and after the D2 dipoles. These components increase the distance from the IP to the first quadrupole Q1 to nearly 53m. The 100mm aperture of the quadrupoles is as large as possible without the coils touching in these twin aperture magnets; the center-to-center distance is 194mm. The reduced number of long-range interactions and the possibility of independent correction systems for the two beams are clear advantages of this design. The major challenges will be designing twin aperture dipoles and quadrupoles with the required field quality, and dealing with the power deposited in the dipoles.

Three alternate designs are shown in Figure 3. The first of these places quadrupoles as early as possible after the beams are separated. This has the advantage of a lower $\beta^\text{max}$ compared to the dipoles first design, and fewer long-range interactions compared to the quadrupoles first design shown in Figure 2. The major challenge will be to design novel twin aperture dipoles and quadrupoles with non-parallel axes. The last two layouts in Figure 3 show very large crossing angles $\sim$8 mrad – in the event that such large crossing angles help increase the luminosity while operating at the beam-beam limit [4].

Table 1 shows the basic IR parameters of the baseline optics and the five options for the upgrade.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Baseline</th>
<th>Fig. 2a</th>
<th>Fig. 2b</th>
<th>Fig. 3a</th>
<th>Fig. 3b</th>
<th>Fig. 3c</th>
</tr>
</thead>
<tbody>
<tr>
<td>IP to Q1 (m)</td>
<td>23</td>
<td>23</td>
<td>52.8</td>
<td>42.5</td>
<td>34</td>
<td>23</td>
</tr>
<tr>
<td>$D_{\text{quad}}$ (mm)</td>
<td>70</td>
<td>110</td>
<td>100</td>
<td>100</td>
<td>100</td>
<td>100</td>
</tr>
<tr>
<td>$\beta^\text{min}$ (cm)</td>
<td>50</td>
<td>16</td>
<td>26</td>
<td>19</td>
<td>15</td>
<td>10</td>
</tr>
<tr>
<td>$\beta^\text{max}$ (km)</td>
<td>5</td>
<td>15</td>
<td>23</td>
<td>23</td>
<td>23</td>
<td>23</td>
</tr>
<tr>
<td>$B_{D1}$ (T)</td>
<td>2.75</td>
<td>15.3</td>
<td>15</td>
<td>14.6</td>
<td>14.5</td>
<td>14.3</td>
</tr>
<tr>
<td>$L_{D1}$ (m)</td>
<td>9.45</td>
<td>1.5</td>
<td>10</td>
<td>12</td>
<td>6</td>
<td>9</td>
</tr>
<tr>
<td>$D_{D1}$ (mm)</td>
<td>80</td>
<td>110</td>
<td>135</td>
<td>165</td>
<td>75</td>
<td>105</td>
</tr>
</tbody>
</table>

Figure 3: Alternate designs for an IR upgrade. Top (a): IR with quads between the separation dipoles, Middle (b): Dipole-first IR with large crossing angle. Bottom (c): Quadrupole-first IR with large crossing angle.

Table 1: IR Parameters
An IR upgrade will be feasible only if the enormous amount of energy deposited in the IR can be safely absorbed. The optics design determines the distribution of the debris power along the IR. In the quadrupole first design, the energy deposited in a magnet increases with the quadrupole length and decreases with aperture. Figure 4, taken from [5], shows the peak energy deposited and the dynamic heat load, as calculated with the MARS code [6], along the four quadrupoles of the triplet at different radial distances for a luminosity of $2.5 \times 10^{34} \text{ cm}^{-2} \text{sec}^{-1}$. The case shown uses a 90 mm aperture quadrupole design (Fig 5.) At a luminosity of $10^{35} \text{ cm}^{-2} \text{sec}^{-1}$, the maximum energy deposited would exceed 4 mW/g in Q2B, well above the quench limit. The heat load to the cryogenics would be greater than 120W/m and a total of 1.6kW would be deposited in the triplet.

With this radiation dose the expected lifetime of G11CR, if used for spacing coils at the magnet ends, would be less than 6 months. Extensive R&D is required for more radiation hard materials. The dynamic heat load is largest in Q1 with electromagnetic showers contributing about 90% of this amount. The dynamic heat load could be reduced with a separate cooling system for an internal absorber inside the Q1.

---

**Figure 4:** Energy deposition in the triplet quadrupoles at a luminosity $2.5 \times 10^{34} \text{ cm}^{-2} \text{sec}^{-1}$, taken from Reference [5]. Top: peak power deposited, Bottom: dynamic heat load

**Figure 5:** Optimized cross-section of a 90 mm aperture Nb$_3$Sn quadrupole, taken from Reference [7].

A preliminary design of large aperture quadrupoles using Nb$_3$Sn with the required field gradient was discussed in Reference [7]. A sketch of the cross-section of a magnet with 2 layer coils and 90mm aperture is seen in Figure 5. This design achieves a field gradient of 205 T/m at a current of 14.1 kA operating at 1.95K. The temperature margin is 2-3 times greater than with NbTi. A major challenge will be to efficiently remove the heat deposited by the beam from the magnet. The eight large wedge shaped holes seen in the cross-section are used to transport superfluid He through the length of the magnet to He tanks outside. Magnets with larger apertures of 100-110 mm would require the use of 4 layer superconductor coils[8].
are several additional design challenges to be addressed for these magnets [7]. For example, the stresses due to Lorentz forces exceed 100 MPa, which will require a strong support structure for the coils.

IV. Dipoles First Design Issues

Dipoles are very efficient at sweeping the charged particle debris from the IP into the magnet due to the large on-axis field. Consequently, the energy deposition issues are more severe for the dipole first IR. Figure 6 shows the distribution of energy deposited on the cross-section of two different magnets: a standard cos θ design and a novel open mid-plane design.

The peak power density with a horizontal crossing angle is in the horizontal plane and skewed in the direction of the outgoing beam. At a luminosity of $10^{35}$ cm$^{-2}$sec$^{-1}$, the peak power density in the horizontal plane in the cos θ design is 50mW/g, about two orders of magnitude larger than the peak power density in the baseline optics. Out of the total power of 9kW from the IP, about 3.5 kW is deposited in the first dipole. The need to minimize the beam power in the superconducting coils has spurred the design of an open mid-plane dipole magnet [10]. Recent refinements have further reduced the amount of energy deposited in the coils [11].

Figure 7 shows particle tracks in a possible layout with an open mid-plane dipole D1 first and absorbers TAS and TAN on either side. The tracks shown originate from a single pp event at the IP and only particles with energies > 10 GeV are shown. TAS just before the D1 absorbs the majority of these particles, and only the most energetic particles propagate to the back end of D1 and the TAN absorber downstream of it. Tungsten rods at liquid nitrogen temperatures are placed in the mid-plane to absorb much of this radiation [11].

Figure 8 shows the power distribution with this design. The peak power density in the coils is significantly reduced by about two orders of magnitude from the cos θ design. However the magnet design itself is quite complex and considerable R&D is required to prove that such a magnet can be built.

V. NbTi Magnets for IR upgrades

It is widely accepted that the road to building accelerator quality magnets with Nb$_3$Sn will be long and hard. It therefore makes sense to
We have seen that a quadrupole first design with Nb₃Sn magnets of 110mm aperture and ~6 m length can achieve $\beta^* = 16$ cm. To achieve the same $\beta^*$ with NbTi magnets which have a lower pole tip field would require magnets with lengths in the range 8-9m to have the same focusing strength. The increased length of magnets implies that $\beta_{\text{max}}$ will be about 30% larger requiring apertures in the range 120-130 mm. Beams are separated further away from the IP implying an increase in the number of parasitic interactions by 15-20%. These accelerator physics issues could possibly be addressed by better correction systems. The smaller temperature margin of NbTi implies that these magnets will be very sensitive to beam heating. This problem could be addressed by placing adequate absorbers within these magnets at the cost of further increasing the aperture and therefore also the length and $\beta_{\text{max}}$.

Radiation hard materials that can be used with NbTi also need to be developed to handle the increased radiation dose. This brief discussion suggests that NbTi quadrupoles may be suitable for a modest luminosity upgrade but probably not for the ultimate upgrade.

VI. Open questions, issues and challenges

The basic IR design concepts discussed here show that $\beta^*$ can be reduced by factors of 2-5 with respect to the baseline design. However this alone does not increase the luminosity by the same factors. At smaller $\beta^*$, the crossing angles have to be increased to keep the same beam separations, which limits the increase in luminosity. Several options to recover this luminosity loss have been considered. Bunches could be shortened with higher frequency rf cavities; increased voltage with more cavities is likely not possible due to the lack of space. However even this option is expensive. Another option is to introduce crab cavities that make the beams collide head-on at the IPs. The estimated voltages are rather large, around 40 MV, and special care will be needed to keep cavity errors to a minimum to prevent emittance blow up.
KEK plans to install crab cavities in their B factory in late 2005 – much will be learned from this first experience with crab cavities in an operating machine. Finally one can also envisage increasing the beam current to the extent allowed by the injectors and potential instabilities in the LHC.

We list here some other accelerator physics questions specific to the IR optics design.

- Can IR magnet errors be adequately corrected given the very large $\beta$ functions?
- Are the very large crossing angle schemes in any way feasible?
- Can dispersion suppressors be designed for quadrupoles with non-parallel axes?

Several magnet R&D challenges have to be met as well. All designs put a premium on achieving very high fields. In the case of quadrupoles first, this maximizes the aperture for a given gradient, while for dipoles this separates the beams quickly and brings quadrupoles in closer to the IP. The operating field is increased from 8T in the baseline optics to 13-15 T in the dipoles or at the pole tip of the quadrupoles. This field can only be achieved today with Nb$_3$Sn technology. All designs also put a premium on large apertures to allow a smaller $\beta^*$. Quadrupole apertures as large as 110 mm may be required, while the need to accommodate both beams in the first dipole D1 require an aperture around 130 mm. The magnet design challenges with these large apertures have to be addressed. Perhaps the hardest will be to meet the various demands of coping with the immense energy deposited: issues of quench stability, heat load on the cryogenics, and developing radiation hard materials have to be addressed.

The first few years of operating the LHC at the nominal luminosity will help to determine the main factors that limit luminosity and will be essential in guiding the design of an upgrade. For example, if active beam-beam compensation is shown to work then the simpler quadrupole first design may be the more attractive option. The requirements of the experiments may also change for the upgrade. They may allow the magnets to move in closer to the IP in the quest for higher luminosity or alternatively the magnets may have to be pushed back if the Higgs can only be found in channels such as WZ scattering [3].

**VII. Summary**

We have reviewed some options for increasing the luminosity with an IR upgrade. There are two “simple” upgrade options with either quadrupoles first or dipoles first using Nb$_3$Sn technology that have the potential of reducing $\beta^*$ by factors of 2-3. More “exotic” layouts have also been presented that might reduce $\beta^*$ by up to a factor of 5. Energy deposition and radiation hardness will be the major challenges to address at $10^{35}$ cm$^{-2}$ sec$^{-1}$ luminosity. While a modest upgrade may be possible using magnets with NbTi technology, the ultimate luminosity goal seems feasible only with Nb$_3$Sn technology. However considerable R&D is required before this new technology is proven. Given the complex and disparate challenges on the road to higher luminosity, all promising options need to be pursued now to ensure success.

**REFERENCES**

FIRST RESULT OF INDUCTION ACCELERATION IN THE KEK PROTON SYNCHROTRON

Ken Takayama1,2,3, K.Koseki2, K.Torikai1,4, A.Tokuchi5, E.Nakamura1,2, Y.Arakida1, Y.Shimosaki1, M.Wake1, T.Kono1, K.Horioka1, S.Igarashi1, T.Iwashita1, A.Kawasaki5, J.Kishiro1,6, M.Sakuda7, H.Sato1, M.Shiho3,6, M.Shirakata1, T.Sueno1, T.Toyama1, M.Watanabe6, and I.Yamane1

1High Energy Accelerator Research Organization(KEK), 1-1 Oho, Tsukuba, Ibaraki 305-0801, Japan
2The Graduate University for Advanced Studies, Hayama, Miura, Kanagawa 240-0193, Japan
3Tokyo Institute of Technology, 2-12-1 Ookayama, Meguro, Tokyo 152-8550, Japan
4Graduate School of Engineering, Kyushu University, 6-10-1 Hakozaki, Fukuoka, Fukuoka 812-8581, Japan
5Nichicon (Kusatsu) Corporation, 2-3-1 Yagura, Kusatsu, Shiga 525-0053, Japan
6Japan Atomic Energy Research Institute, 2-4 Shirane, Tokai, Naka, Ibaraki 319-1195, Japan
7Physics Department, Okayama University, 3-1-1 Tsushimanaka, Okayama, Okayama 700-8530 Japan

Abstract

Induction acceleration of a single RF bunch was confirmed in the KEK PS.

INTRODUCTION

Four years ago, the concept of an induction synchrotron employing induction accelerating devices was proposed by Takayama and Kishiro [1] for the purpose to overcome the shortcomings, such as a limitation of the longitudinal phase-space available for acceleration of charged particles, in an RF synchrotron, which has been one of indispensable instruments for nuclear physics and high energy physics since the invention by McMillan [2] and Veksler [3]. Accelerating devices in a conventional synchrotron, such as an RF cavity, are replaced with induction devices in the induction synchrotron. A gradient focusing force, seen in the RF waves, is not indispensable for the longitudinal confinement of particles. Pulse voltages, which are generated at both edges of some time-period with opposite sign, as shown in Figure 1, are capable of providing longitudinal focusing forces. A pair of barrier-voltage pulses work in a similar way to the RF barrier, which has been demonstrated at FNAL and BNL [4]. The acceleration and longitudinal confinement of charged particles are independently achieved with induction step-voltages in the induction synchrotron. This notable property of the separated-function, in the longitudinal direction brings about a significant freedom of beam handling never seen in a conventional RF synchrotron, in which radio-frequency waves in a resonant cavity simultaneously take both roles of acceleration and longitudinal confinement. A big difference in the phase space between an RF synchrotron and the induction synchrotron is schematically shown in Figure 2.

FIGURE 1. Conceptual view of the induction synchrotron and super-bunch with barrier voltages and acceleration voltage

FIGURE 2. RF bunches and Super-bunch in the phase space

Associated with the separated-function, various figure of merits are expected. The formation of a super-bunch, which is an extremely long-bunch with a uniform line-density, and the use of which is considered in a proton driver for the second-generation of neutrino physics and a
future hadron collider [5], is most attractive. In addition, transition crossing without any longitudinal focusing seems to be feasible [1], which could substantially mitigate undesired phenomena, such as bunch shortening due to non-adiabatic motion and microwave instabilities [6].

In April 2003, we started a new project to demonstrate an induction synchrotron using the KEK PS, where a super-bunch will be accelerated with the induction devices. The project consists of three stages [7]: at the first stage a single bunch captured in the RF bucket is accelerated up to 8GeV with the induction accelerating system, at the second stage multi RF bunches injected from the 500MeV Booster are captured in the barrier bucket to merge into a single super-bunch, and the last stage this super-bunch will be accelerated up to the flat-top energy. Here the current status of the first stage is presented.

The time-duration between barrier-voltage pulses determines a size of the super-bunch. For acceleration of the super-bunch, a long accelerating step voltage with a pulse length of the order of µ-sec is required. The voltage has to be generated at the revolution frequency of the beam in a ring. In the case of a ring-circumference of 300 m, a repetition rate of 1 MHz is required. So far, there has been no induction accelerating system capable of meeting these parameters. Recently, the prototype devices, which can generate a 250nsec flat-top voltage at a repetition rate of cw 1 MHz, have been assembled at KEK after the 3 years R&D stage and combined with the existing RF accelerating system.

A single proton bunch trapped in an RF bucket was accelerated with the induction accelerating system from 500MeV to about 1.5 GeV. In this paper, we report on the first experimental result of induction acceleration of a single RF bunch in the KEK proton synchrotron (KEK-PS) as well as a brief description of the induction accelerating system.

INDUCTION ACCELERATION SYSTEM

The key devices required to realize an induction synchrotron are an induction accelerating cavity [8] and a pulse modulator driving the cavity [9]. These devices are notably different from similar devices employed in modern linear induction accelerators [10]. Remarkable properties are its switching characteristics, repetition ratio, and duty factor. Another different feature is that the pulse modulator has to be kept far from the induction cavity placed in the accelerator tunnel, because the solid-state power switching elements obtainable at present can’t survive an extremely high radiation dose. Thus, the pulse modulator is connected with the accelerating cavity through a long transmission cable. In order to reduce reflection from the load, a matching resistance has been installed at the end of the transmission cable. The induction accelerating system consists of an induction cavity with a matching load, a transmission cable, a pulse modulator, and a DC power supply. A typical system capable of generating a step-pulse of 2 kV output voltage and 18 A output peak current at 1 MHz with 50% duty has been demonstrated at KEK. An equivalent circuit for the induction accelerating system is shown in Fig.ure 3.

FIGURE 3. Equivalent circuit for the induction acceleration system

A core material of the induction cavity employed for the first acceleration experiment is a nanocrystalline alloy, called Finemet (Hitachi Metal). Heat generated due to core loss is cooled down by insulation oil. The electrical parameters of each unit-cell, such as capacitance, inductance, and resistance, which determine the property of pulse rising and falling are listed in Table 1. Three unit-cells with a 2 kV output voltage per unit are mechanically combined into a single module for the convenience of installation. Since the inner conductor with three ceramic gaps is common to three unit-cells, but both sides of each gap are electrically connected to the outer edges of each cell, a particle is accelerated with the same voltage passing these gaps. An engineering drawing is depicted in Figure 4.

<table>
<thead>
<tr>
<th>Capacitance (pF)</th>
<th>260</th>
</tr>
</thead>
<tbody>
<tr>
<td>Resistance (Ω)</td>
<td>330</td>
</tr>
<tr>
<td>Inductance (H)</td>
<td>110</td>
</tr>
</tbody>
</table>
A full-bridge switching circuit in the pulse-modulator, which is depicted in Figure 5, was employed because of its simplicity. The pulse modulator is capable of generating bipolar rectangular shaped voltage pulses. The full-bridge type pulse-modulator consists of four identical switching arms. Each switching arm is composed of 7 MOS-FETs, which are arranged in series. Their gates are driven by their own gate-driving circuits. The gate signals are generated by converting light signals provided from the pulse controller, which is a part of the accelerator control system, to electronic signals. Its details will be presented elsewhere [10]. The inside figure is given in Figure 6.

**FIGURE 4.** Induction cavity

**FIGURE 5.** Architecture of the power modulator and switching sequence (top), typical pulse pattern (bottom)

**FIGURE 6.** Photo of the pulse modulator

HYBRID ACCELERATION

The entire system employed for the current experiment is schematically shown in Figure 7. The generation of a 2 kV voltage pulse is directly controlled by trigger pulses for the switching elements of the pulse modulator, a master signal of which is created in the digital signal processor (DSP) synchronizing with ramping of bending magnets, and the gate-driving signal patterns initiated by this master trigger-signal are generated by the following signal-pattern generator to be sent to the gate controller of the pulse modulators through a long coaxial cable. The DSP counts the B-clock signal, and achieves the desired revolution frequency. Any delay between the accelerating pulse and a bunch monitor signal is always corrected by the DSP. The system is connected to the existing RF system through the RF signal, which shares the B-clock signal. As a result, synchronized induction acceleration is guaranteed. Here, the RF does not contribute to acceleration of the beam bunch but gives the focusing force in the longitudinal direction; the beam bunch is in principle trapped around the phase of zero. The machine parameters of the KEK PS employed for the experiment are listed in Table 2.
FIGURE 7. Schematic view of the hybrid accelerating system of the KEK PS

Table 2

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Circumference</td>
<td>$C_0$ 339 m</td>
</tr>
<tr>
<td>Transition energy</td>
<td>$\gamma_t$ 6.63</td>
</tr>
<tr>
<td>Injection/extraction energy</td>
<td>500 MeV/8 GeV</td>
</tr>
<tr>
<td>Revolution frequency</td>
<td>$f_0$ 668 – 877 kHz</td>
</tr>
<tr>
<td>Ramping time (transient for start/stop)</td>
<td>1.9 sec (100 msec)</td>
</tr>
<tr>
<td>RF voltage</td>
<td>$V_{rf}$ 48 kV</td>
</tr>
<tr>
<td>Harmonic number</td>
<td>$h$ 9</td>
</tr>
<tr>
<td>Induction voltage per turn</td>
<td>$V_{ind}$ 5.2 kV</td>
</tr>
</tbody>
</table>

**EXPERIMENTAL RESULTS**

In the experiment, the signals of the bunch monitor and three current transformers (CT), which always observed the current flow through the matching resistances, were monitored on the digital oscilloscope located at the central accelerator control room (CCR). Before the experiment, an actual induced voltage at the ceramic gap, an output voltage of the transmission cable, and the CT signal were simultaneously measured and the correspondence between each other was well calibrated. In addition, a delayed timing of the master gate signal triggering the pulse modulator is adjusted by the DSP so that the bunch signal would stay around the center of the induction voltage pulse through the entire accelerating period. Typical wave-forms of the CT signals are shown in Figure 8 together with the bunch signal.

**THEORETICAL BACKGROUND**

Under coexistence with the RF voltage and the induction voltage, a charged particle receives an energy gain per turn,

$$eV_{acc}(t) = e[V_{rf} \sin \phi(t) + V_{ind}]$$  \hspace{1cm} (1)

where $V_{rf}$ and $V_{ind}$ are the RF voltage and the induction voltage, respectively, and $\phi(t)$ is the position of the particle in the RF phase, $\alpha_0$. The orbit and energy of a particle are dominated by the following equations: for the force balance in the radial direction

$$m\gamma \cdot \left( \frac{e\beta}{\rho} \right)^2 = e\epsilon \cdot B(t)$$  \hspace{1cm} (2)

where $B(t)$ is the bending field and $\rho$ is the bending radius, and for the change in energy

$$mc^2 \cdot \frac{d\gamma}{dt} = e\epsilon \cdot \frac{C_0}{C_0} \cdot V_{acc}(t)$$  \hspace{1cm} (3)

From (2) and (3), the accelerating voltage must satisfy the relationship, $V_{acc}(t) = \rho \cdot C_0 \cdot dB/dt$, so that the particle is synchronously accelerated with ramping of the bending field. The bending field is linearly ramped over 1.7 sec. In this linear ramping region, $V_{acc}$ of 4.7 kV is required. For the simplicity, the induction voltage was fixed to be close to 5.1 kV.

**EXPERIMENTAL OBSERVATION**

In order to confirm the induction acceleration, the phase signal, which shows the relative position of the bunch center to the RF, was measured through an accelerating region of concern. Particularly, we focused on three cases:
(1) with an RF voltage alone, (2) with an RF voltage and a positive induction voltage, and (3) with an RF voltage and a negative induction voltage. From Eq.(1), a theoretical prediction is
\[ \phi_s = \sin^{-1}\left(\frac{V_{acc}}{V_{rf}}\right) \]
~5.7 degrees for case (1), \[ \phi_s = -0.66 \text{ degree} \] for case (2), \[ \phi_s = \sin^{-1}\left(2\frac{V_{acc}}{V_{rf}}\right) \] ~12.4 degrees for case (3), where \( \phi_s \) represents the position of the bunch center in the RF phase. Since the induction voltage is devoted to the acceleration for case (2), the RF does not serve for the acceleration, but takes a role of capturing alone; thus, the phase must be zero. In case (3), the RF has to give a two-times larger energy to the bunch than case (1) from the energy-conservation law; the phase should increase by a factor of two. Actually, the time-evolution of the phases through acceleration has been observed, as seen in Figure 9. At a first glance, we can find the qualitative agreement with the theoretical prediction. For both side of the transition, the evolution in the phase difference is clearly understandable.

This result was obtained at the very preliminary stage of the experiment. More recent results and their better understanding are available in the paper [11].

**ACKNOWLEDGEMENT**

The present research has been supported by the KEK PS division. The authors acknowledge S. Ninomiya for providing a lot of information on the KEK-PS RF system and helpful comments on the experimental result, and D. Arakawa for an advice on the bunch monitor. The present research has been financially supported by a Grant-In-Aid for Scientific Research for Creative Scientific Research (KAKENHI 15GS0217), and its early stage was partially supported by a Grant-In-Aid for Scientific Research on Priority Area (KAKENHI 140646221).

**REFERENCES**


http://www.arxiv.org/pdf/physics/0412006

**SUMMARY**

It is emphasized that for the first time charged particles in the high-energy accelerator ring were accelerated with the induction accelerating system. The experimental fact of the induction acceleration in a circular ring and the reality that the key devices developed for this purpose, such as a pulse modulator as a switching driver, are on our hands are big mile-stones for us to achieve an induction synchrotron in near future and a super-bunch hadron collider in no so far future. Last it is emphasized that the induction system worked well through the entire operating period of 24 hours without any trouble. We will conclude that a typical demonstration will be performed for actual applications.

**FIGURE 9.** Temporal evolution of the phase
Abstract

It is preferred to make a larger crossing angle to avoid parasitic collisions between beams in LHC. The transverse-longitudinal aspect ratio of the LHC beam is 17 $\mu$m/7.5 cm = 0.23 $\times$ 10^{-3}. Crossing angle larger than the aspect ratio causes geometrical degradation of luminosity.

Crab cavity can realize head-on collision without crossing angle effectively, with the result that it recovers the luminosity performance. However crab cavity can also be a source of diffusion: that is, Jitter of RF phase of crab cavity and main acceleration cavity induce a random transverse offset at the collision point, which causes diffusion due to beam-beam interaction.

We discuss feasibility of the crab cavity for hadron colliders from the viewpoint of dynamics of the beam-beam interaction.

INTRODUCTION

A crab cavity has been designed and constructed at KEK to boost up the luminosity of KEK-B factory (KEKB). In collision with a crossing angle in very high intensity, Arnold diffusion make worse the luminosity performance. Crab cavity, which realizes head-on collision effectively, is expected to make increase the beam-beam parameter twice [1].

The crab cavity generates z-dependent dispersion, $\Delta z = z_{cr} \Delta z$, which cancels the dispersion induced by crossing angle effectively. The transfer matrix of arc, crab cavity and crossing effect is block diagonalized to 2 $\times$ 2 matrices, except for small nonlinear kinematical terms, with the result that the luminosity performance is improved. This scenario is for $e^+e^-$ colliders with very high beam-beam parameter, $\xi \approx 0.1$.

In LHC, the transverse-longitudinal aspect ratio of the beam is 17 $\mu$m/7.5 cm = 0.23 $\times$ 10^{-3}. Therefore crossing angle more than 0.2-0.3 mrad causes geometrical loss of luminosity. For such small crossing angle, parasitic collision between beams is serious [2]. Crab cavity manages both week parasitic interaction and no geometrical loss.

In the case of LHC, the beam-beam parameter is much less than that of $e^+e^-$ colliders, therefore the diffusion is much weaker. However since there is no damping mechanism due to the absence of synchrotron radiation, even a slow diffusion can be serious for the hadron colliders. One important subject is how strong Arnold diffusion in LHC. Crab cavity works to improve the luminosity performance from this point of view, if the diffusion is dominant for the beam-beam effect in LHC.

Crab cavity is operated with a transverse mode like TM110, and deflected a bunch on $x-z$ plane by choosing rf phase. Voltage required for a deflection angle $\phi$ is expressed by

$$ V = \frac{eE \tan \phi}{\omega_{RF} \sqrt{\beta_{x,\omega} / \beta_{x}^{*}}} $$

To realized collision without crossing angle effectively, the deflection angle should be equal to the half crossing angle.

The crab cavity can be source of diffusion: i.e., since the crab cavity is operated by a transverse mode, the deviation and jitter of RF phase give a dipole kick to the beam, with the result that transverse offset at the collision point is generated. Both of phases of main RF and crab cavity can be source of the transverse offset.

Jitters of RF phase of main cavity causes a deviation of timing of beam arrival at the crab cavity. The transverse offset, which arise from the jitter of main RF system, is expressed by

$$ \delta x = \frac{e \tan \phi}{\omega_{RF}} \delta \psi_{RF} $$

where $\delta \psi_{RF}$ is the phase error of the main RF system.

The crab cavity gives a transverse kick due to its jitters of RF phase, with the result that the offset given by the kick is expressed as follows,

$$ \delta x = \frac{e \tan \phi \cos (\pi \nu_{RF} - \Delta \psi(s^*, \xi_{cr}) \delta \psi_{cr,RF}}{2 \sin \pi \nu_{RF}} $$

where $\Delta \psi(s^*, \xi_{cr})$ and $\delta \psi_{cr,RF}$ are the betatron phase difference between the collision point and the crab cavity and the deviation of the RF phase of the crab cavity, respectively. In the both cases, the jitter of transverse offset is given by $\delta x \approx e \tan \phi \delta \psi_{cr,RF}$.

In hadron colliders, the beam-beam limit seems to be determined by diffusion. Therefore it is important to understand which diffusion mechanism is dominant. Crab cavity reduces Arnold diffusion, while diffusion due to phase errors is introduced. Many other diffusion sources seem to exist in LHC, and some of them affect the beam-beam performance.

We study the effect of beam-beam performance of crab cavity. Parameter of LHC is shown in Table 1.

PARAMETER OF THE CRAB CAVITY

A crab cavity with squashed cell has been developed to install to KEKB. The crab cavity generates V = 1.44 MV. The crab cavity is designed for B factory to reduce...
impedance of parasitic modes including lowest mode, because the operating current is very high, 2.6 A for the design, and is increased to 10 A future.

The half crossing angle is 0.15 mrad at an early stage of LHC, and is proposed to increase to 4 mrad to reduce the effect of parasitic collisions. Required voltage is evaluated to be $V_{crab}=15$ MV/mrad. The voltages are 2.25 MV for 0.15 mrad and 60 MV for 4 mrad, and the number of cavity is 4 and 100 per ring, respectively. The number is feasible for 0.15 mrad, but is questionable for 4 mrad. We may need another choice, for example multi-cell type, for such a large crossing angle.

The beam intensity of LHC ($I=600$ mA) is comparable with that of $e^+e^-$ ring, though energy is much higher. Care for coupled bunch instability due to parasitic cavity mode may be necessary. KEKB type of crab cavity is investigated for LHC as a first trial, though RF frequency 500 MHz is not match for LHC (400 MHz).

$$\tau_{G,L} = \frac{MN\gamma R}{4\pi \gamma T_0} \omega R(Z(\omega))_{L,peak}$$

$$\tau_{G,T} = \frac{MN\gamma R}{4\pi \gamma T_0} \omega R(Z(\omega))_{L,peak}$$

The impedance of parasitic modes have been estimated at KEKB using MAFIA by Akai et al. [3]. The peak values of impedance and corresponding growth rate were given as

$$\omega Z(\omega)_{L,peak} = 13k\Omega \cdot G Hz \quad \tau_{G} = 1.5 \text{sec}$$

$$Z(\omega)_{T,peak} = 0.025 M\Omega/m \quad \tau_{G} = 1.5 \text{sec}$$

The impedance and growth time is for one cavity, the value should be devided by 4 or 100 for the crossing angle, 0.15-4 mrad. Therefore the growth rate is 0.4 s or 15 msec for 0.15 or 4 mrad, respectively.

If we use multi-cell cavity, the number is reduced, while parasitic modes may be more serious.

**BEAM-BEAM EFFECT**

Jitters of RF phase of crab cavity and/or main cavity induce a transverse offset at the collision point. The noise of collision offset induces diffusion due to the beam-beam interaction [4], with the result that emittance growth occurs.

The emittance growth is evaluated by a strong-strong simulation code named by BBSS [7]. 100,000 macro-particles are used to represent a bunch and a bunch is divided into 5 slices. Collisions and revolutions are repeated during 10,000 turns, and beam size and luminosity were estimated turn by turn. These values, the number of macro-particles and turns, are sufficient to study to survey of beam-beam diffusion briefly, though long term stability can not be discussed accurately.

First of all, we show the behavior of luminosity and beam size without phase error in Figure 1. Their evolutions with or without crab cavity are depicted for half crossing angle of 0.15 mrad. The beam size and luminosity were kept constant during 10,000 turns as shown in the figure. Arnold diffusion due to the crossing angle was not seen for the statistics of the macro-particles and the number of revolutions. Since the beam is tilt due to crab cavity on x-z plane, its size is enlarged as a matter of form. Though beam size for crab crossing is larger than that without it, the luminosity is also larger than that without it.

![Figure 1: Evolution of Luminosity and horizontal beam size for no crab cavity](image)

We next discuss the fluctuation of the transverse offset due to the phase jitter of RF system. Fluctuation of trans-
verse offset was included using Gaussian random number in the simulation. If the noise of the jitter has a pure white spectrum, the correlation function of the fluctuation is \( \langle x(n) x(m) \rangle = \delta x \delta \tau_{nm} \). The fluctuation with correlation time, \( \tau \), \( \langle x(n) x(m) \rangle = \delta x e^{-|n-m|/\tau} \). The fluctuations with the correlation time are generated by random numbers [6].

\[
x(n + 1) = x(n) + \delta x / \sqrt{1/\tau_{CG}} \tag{8}
\]

where \( \tau_{CG} \) is Gaussian random number with unit deviation.

Diffusion due to the phase jitter of accelerating cavity is now discussed. A bunch is kicked by the jitter of arrival time of bunches. In the simulation transverse offset is given via longitudinal displacement: i.e., \( \delta x \) is not applied, but \( \delta z \) is applied directory. \( \delta x \) is given by \( \delta z \) via \( \zeta_x \) induced by the crab cavity. Figure 2 shows evolution of beam size and luminosity for various error strengths. For largest strength the diffusion is clearly visible, but the diffusion decrease rapidly for decreasing the error strengths. The correlation time of the jitter is assumed to be 100 turns for the time being. The time should be reconsidered with taking into account of the quality factor of the cavity or other parameters.

The diffusion rate is estimated by growth of \( \sigma_x^2 \) per turn in Figure 2. The diffusion rate as a function of \( \delta x \) is shown in Figure 3. As is shown latter in an analytical estimate, the diffusion rate is a quadratic function of \( \delta x \). If we use the dependence, the diffusion rate is approximately represent by

\[
D[\mu m^2] = \langle \Delta \sigma_x^2 \rangle \approx 1.4 \times 10^{-3} |\delta x(\mu m)|^2. \tag{9}
\]

Realistic value of the phase jitter is considered to be 0.1-1 degree, therefore jitter of longitudinal offset and resulting transverse offset are obtained as \( \delta z = 0.2 - 2 \) mm and \( \delta x = 0.03 - 0.3 \) \( \mu \)m, respectively, for 0.15 mrad. The diffusion rate is given as \( 1.26 \times 10^{16} - 10^{18} \) m² per turn, with the result that the luminosity life time is estimated to be 100-10,000 sec \( (10^6 - 10^8 \) turns).

The same calculations were done without crab cavity for crossing angle of zero and 0.15 mrad. In this calculation, a beam experiences jitter of longitudinal position of the other beam at the collision, but does not do transverse jitter. Figure 4 shows the evolution of beam size and luminosity. We can see modulation of luminosity, but can not see its monotonically decrease: that is, this fact indicates that transverse jitter mainly contributes the diffusion.

We next discuss the jitter of the crab cavity phase. In
this case, jitter of transverse offset is applied directly. Figure 5 shows the evolution of luminosity and beam size $\delta x$ is $17 \mu m$ with time correlation of 100. The offset corresponds phase error of 5 degree. Those for RF phase jitter with the same amplitude is depicted as a reference. The luminosity degradation and beam size diffusion are comparable, though those for crab phase are somewhat smaller than those for RF phase.

**Figure 5**: Evolution of Luminosity and horizontal beam size for phase jitter of crab cavity. Those for RF phase jitter with the same $\delta x$ is depicted as a reference.

### ANALYTIC ESTIMATE OF THE DIFFUSION RATE

Collision with transverse offset causes diffusion of motion of beam-particle via tune variation [4]. The diffusion rate $D(J_x)$ defined by fluctuation of $J_x^2$ of beam particle per turn: i.e., $D(J_x) = \langle J_x^2(N) - J_x^2(0) \rangle /N$, where $N$ is turn. The diffusion rate is expressed in one dimensional model by

$$D(J_x) = \frac{C\sigma|\delta x|^2}{8 - 4/\tau} \sum_{k=0}^{\infty} \frac{\sinh \theta(2k + 1)^{1/2}G_k^2(\theta)}{\cos \theta - \cos [2\pi(2k + 1)|\nu|]}. \tag{10}$$

Some functions and variables in Eq.(10) are defined by

$$G_k = \frac{\sqrt{\Delta}}{\sigma} \left[ U'_{k+1} + U'_k \right] - \frac{1}{\sqrt{\Delta \sigma}} \left[ (k + 1)U_{k+1} - kU_k \right], \tag{11}$$

where

$$U_k(\alpha) = \int_0^\alpha \frac{1}{w} \left[ \delta_{0k} - (2 - \delta_{0k})(-1)^k e^{-wJ_k(w)} \right] dw. \tag{12}$$

and

$$\theta = -\ln(1 - 1/\tau) \quad C = \frac{N_p r_p}{\gamma_p} \quad \alpha = \frac{\beta^* J_x}{2\sigma^2} \tag{13}$$

We evaluated $D$ as $D(\alpha = 1) = 1.5 \times 10^{-25} \text{m}^2/\text{turn}$. The simulation showed $D = 10^{-28} \text{m}^2/\text{turn}$. The theory was extended to two dimensional model [5]. We should compare the diffusion rate with the two dimensional model.

### SUMMARY

We studied crab cavity option for LHC. As a first trial, KEKB type of crab cavity was evaluated for LHC, since it is cared for high current operation. Required voltage growth rate of coupled bunch instability caused by parasitic mode of the cavity was estimated to be 0.4 s and 15 msec for the crossing angle of 0.15-4 mrad, respectively.

We studied a degradation of beam-beam performance caused by phase jitter of accelerating and crab cavities. The diffusion rate of the beam size was evaluated to be $D[\mu m^2]$ $\approx 1.4 \times 10^{-31} |\delta x(\mu m)|^2$. If the phase jitter is between 0.1-1 degree, the transverse offset 0.03-0.3 $\mu m$ arise from the jitter, and luminosity life time is estimated as 100-10,000 sec. Time correlation was chosen to be 100 in the evaluation. Shorter correlation time make worse the diffusion.

### ACKNOWLEDGMENTS

The authors acknowledge members of KEKB crab cavity group, especially, K. Akai and K. Hosoyama, for information of the crab cavity design. The author thanks to T. Sen and F. Zimmermann for discussions on effect of noises in the beam-beam interaction.

### REFERENCES

Beam-Beam Compensation Schemes

F. Zimmermann, CERN, Geneva, Switzerland

Abstract

Several techniques have been proposed for compensating the effects of the head-on or long-range beam-beam interaction. These include alternating crossing at several intersection points, electromagnetic wires, electron lenses, and superbunches. I discuss some prominent compensation schemes and their R&D status, with special emphasis on the LHC and its upgrade.

1 HISTORY AND PERSPECTIVE

The beam-beam interaction has limited the performance of almost all past and present storage-ring colliders. Therefore, not surprisingly, the history of accelerators has witnessed a number of ingenious attempts to overcome this limitation by various compensation techniques.

Famous are the 4-beam collisions with charge neutralization conceived around 1970 at the Orsay DCI [1]. Without net charges, all forces were cancelled in order to produce a vanishing beam-beam tune shift. Unfortunately, small offsets in the centroid positions of the oppositely charged beams were amplified exponentially, and the DCI performance was limited by the resulting collective beam-beam instabilities [2].

In 1975, head-on collisions with a high beam-beam tune shift were simulated in the CERN ISR by two vertically separated copper bars carrying 1000-A of electric current [3]. The ISR experiment showed that resonances of order 10 or higher contributed to the proton-beam-beam limit. Although this was not strictly speaking an attempt at compensation, it pointed the way to future compensation experiments, like those employing current-carrying wires (see below).

In 1982, the use of octupoles for reducing the beam-beam tune spread was studied for LEP by means of simulations [4]. Later, octupoles were indeed shown to be beneficial, especially for background control, at BEPP-4, VEPP-2M [5, 6] and DAFNE [7].

Plasma-based compensation schemes were theoretically explored both for linear electron-positron [8] and for muon colliders [9], but have not yet been tested in practice.

Compensation of beam-beam effects in hadron colliders by a low-energy electron beam was proposed for the SSC [10] and the LHC [11]. Since 2001, the first practical realization of such a scheme, the Tevatron Electron Lens (TEL), [12] is operational at FNAL. The TEL collides the antiprotons with an electron beam of appropriate shape and strength to compensate for the effect of the antiproton-proton collisions at other locations of the ring. The TEL can be pulsed with a different electron current for each bunch. In the near future, a pair of TELs will be capable of reducing both the intra- and inter-bunch tune spread in the Tevatron antiproton beam [12, 13].

The present and future generations of hadron colliders (Tevatron Run-II, LHC, RHIC upgrade) accommodate not only head-on collisions at the design ‘interaction points’ (IPs) inside the particle-physics detectors, but also a significant number of ‘long-range collisions’ (sometimes called ‘parasitic collisions’), where the beams are separated by a few times the transverse rms beam size, but still notice each other’s fields.

As a concrete example, Fig. 1 shows the overall LHC layout, with four primary interaction points (IPs), at the experimental detectors of ATLAS, ALICE, CMS, and LHC-B, respectively. On the way to and from each primary IP, the LHC proton bunches suffer a maximum of 30 long-range collisions (15 on each side of the IP), as is illustrated in Fig. 2. Summed over the entire ring, the total number of long-range collisions experienced by a regular proton bunch is 120.

The LHC beam consists of 39 trains of 72 bunches each, spaced by 25 ns. The trains are separated by gaps of 574, 599 or 2595 ns, in order to accommodate the rise and fall times of various injections and extraction kickers in the LHC and its injectors. Due to these gaps, where bunches in the opposing beam are missing, proton bunches at the end or start of a bunch train do not experience the full number of long-range collisions, and, as a result, may have different orbits and tunes. These bunches could exhibit a reduced lifetime. They are therefore called ‘PACMAN’ bunches. Only about half of the LHC bunches are regular bunches, all others belong to one of many different equivalence classes of PACMAN type [15].

The long-range collisions are a concern for all bunches, because they perturb the motion of protons at large betatron amplitudes, where they come close to the opposing beam. Thereby, they generate a ‘diffusive aperture’ [16], beyond which a particle is rapidly lost, i.e., the diffusive aperture can be considered equal to the short-term dynamic aperture. The diffusive aperture induced by the long-range collisions may result in high background at the experiments and in a poor beam lifetime. The effect of long-range collisions is a problem of increasing importance, from the SPS over the Tevatron Run-II to the LHC, i.e., for operation with a larger number of bunches, as is illustrated by Table 1.

The experience from Tevatron Run-II confirms the potential danger of long-range encounters. Quoting T. Sen [17], ‘Long-range beam-beam interactions in Run II at the Tevatron are the dominant sources of beam loss and lifetime limitations of anti-protons, especially at injection en-
Table 1: Number of long-range collisions in some hadron colliders.

<table>
<thead>
<tr>
<th>Collider</th>
<th>No. of Long-Range Encounters</th>
</tr>
</thead>
<tbody>
<tr>
<td>SPS</td>
<td>9</td>
</tr>
<tr>
<td>Tevatron Run-II</td>
<td>70</td>
</tr>
<tr>
<td>LHC</td>
<td>120</td>
</tr>
</tbody>
</table>

points, and in Section 3 the use of electron lenses. Section 4 discusses electro-magnetic wires, Section 5 super-bunches, and Section 6 crab cavities. The review is summarized in Section 7.

2 ALTERNATING CROSSING

2.1 Motivation

The LHC baseline design foresees an alternating crossing scheme, where the two proton beams are crossed horizontally in IPs 1 and 2 (ATLAS and ALICE), and vertically in IPs 5 and 8 (CMS and LHC-B) [25, 15]. The alternating crossing cancels the long-range beam-beam tune shift between the two IPs, and it so assures that PACMAN bunches have (nearly) the same tune as regular bunches. This type of crossing scheme was first proposed by D. Neuffer and S. Peggs for the SSC [18], exactly for this reason, because it reduces the bunch-to-bunch tune spread. The alternating crossing can be considered as a beam-beam compensation scheme, whereby contributions from the beam-beam interaction at one IP are cancelled against those from another IP.

Since, to ensure dynamic vacuum stability, at the LHC a flat beam screen shall be installed in the final-triplet quadrupoles, the orientation of the crossing planes will be frozen for each IP. Therefore, validation of the LHC baseline scheme by simulations and experiments appears prudent. Alternative crossing schemes, while introducing a bunch-to-bunch tune spread, would offer some other advantages, like vanishing vertical dispersion (for $xx$ crossing), easier long-range beam-beam compensation (for $yy$ crossing) or equal background conditions in the two experiments (for $xx$ or $yy$ crossing). In addition, the equal-plane crossing schemes might be intrinsically more stable. A further alternative scheme would be the 'inclined hybrid' crossing, where the two beams are collided in a plane oriented at $45^\circ$ or $135^\circ$ with respect to the horizontal direction [26].

2.2 Simulations

Figure 3 separately illustrates the head-on tune footprint, and the additional tune spreads due to the long-range effects at LHC IP 1 and 5, respectively. Tunes are displayed for particles with transverse starting amplitudes $\vec{r} = (\vec{x}^2 + \vec{y}^2)^{1/2}$ between 0 and 6 times the rms beam size ($6\sigma$). The figure indicates how the alternating cross-
ings in IP1 and IP5 results in a partial cancellation of the long-range tune shifts. Figure 4 compares the total LHC tune spread for a nominal bunch and for an extreme PAC-MAN bunch, i.e., a bunch which only encounters half of the maximum number of long-range collisions. The zero-amplitude tunes of regular and PACMAN bunches are almost the same, thanks to the alternating crossing (the small difference arises due to a lack of symmetry in IP 2 and 8). Note that the total tune spread of the entire LHC beam, including the PACMAN bunches, must fit between harmful resonances in the tune plane, which will limit the maximum achievable tune shift parameter $\zeta$ and thus the bunch intensity $N_b$. This condition is fulfilled more easily with the alternating crossing. Another interesting, and more worrisome, aspect of Fig. 4 is that the long-range collisions ‘fold’ the footprint, i.e., they give rise to a change of direction of detuning with amplitude. This folding is potentially destabilizing, as resonance islands in this amplitude region can be large.

Figure 5 shows simulated diffusion rates for $xy, xx$ and $yy$ crossing in LHC as a function of the starting amplitude $r = x = y$. Here the particles are launched along the diagonal of the $xy$ plane, i.e., the start amplitudes in the $x$ and $y$ planes are equal with randomly assigned initial betatron phases. Results are presented for two different lattice working points, which are mirrored about the tune-plane diagonal (this is the direction of the long-range beam-beam tune shift). In the case of $xy$ crossing, we notice a steep increase in diffusion at an amplitude of about $6.5\sigma$, while for the equal-plane crossing scheme, high diffusion levels are only reached for amplitudes of $9\sigma$. In these simulations, two symmetrically spaced IPs were assumed, the horizon-

Figure 3: Tune footprints due to head-on and long-range beam-beam effects in LHC IPs 1 and 5. Vertical axis refers to the vertical tune, horizontal axis to the horizontal. (Courtesy H. Grote, 2001)

Figure 4: Total tune footprints in the LHC for a regular bunch and for a PACMAN bunch. (Courtesy H. Grote, 2001)

Figure 5: Simulated diffusion rate due to long-range collision with alternating or equal-plane crossing around two LHC IPs at two working points, with fractional tunes mirrored about the diagonal. Shown is the increase in the action variance per turn in units of the square of the nominal emittance, as a function of start amplitude $x = y$ in $\sigma$ [27].

To ascertain that the above simulations of $xy$ crossing and equal-plane crossing were representative, they were re-
peated and extended, this time keeping the zero-amplitude tunes including head-on and long-range beam-beam tune shifts constant, rather than the lattice tunes. We scanned over a large range of vertical tunes to probe the sensitivity to the working point, and we also varied the integer tune by 1 unit to explore the dependence on this parameter. The simulation parameters for this study are compiled in Table 2.

Table 2: Simulation parameters for crossing-plane study; listed values are used if not mentioned otherwise.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Symbol</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>horizontal tune at $A = 0$</td>
<td>$Q_x$</td>
<td>64.30268</td>
</tr>
<tr>
<td>vertical tune at $A = 0$</td>
<td>$Q_y$</td>
<td>59.31268</td>
</tr>
<tr>
<td>bunch population</td>
<td>$N_b$</td>
<td>$1.15 \times 10^{11}$</td>
</tr>
<tr>
<td>beta function at IP</td>
<td>$\beta^{*}_{x,y}$</td>
<td>0.55 m</td>
</tr>
<tr>
<td>relativistic gamma</td>
<td>$\gamma$</td>
<td>7461</td>
</tr>
<tr>
<td>normalized emittance (1σ)</td>
<td>$\gamma_{x,y}$</td>
<td>3.75 μm</td>
</tr>
<tr>
<td>full crossing angle</td>
<td>$\theta_c$</td>
<td>285 μrad</td>
</tr>
<tr>
<td>no. of IPs</td>
<td>$N_{IP}$</td>
<td>2</td>
</tr>
<tr>
<td>no. of parasitic collision points</td>
<td>$N_{par}$</td>
<td>30</td>
</tr>
<tr>
<td>rms beam size at IP</td>
<td>$\sigma^{*}_{x,y}$</td>
<td>16.6 μm</td>
</tr>
<tr>
<td>rms beam divergence at IP</td>
<td>$\theta^{*}_{x,y}$</td>
<td>30.2 μrad</td>
</tr>
</tbody>
</table>

Figure 6 shows the simulation result for the nominal tune. Clearly the $xy$ crossing exhibits the smallest diffusive aperture of about $6\sigma$, independent of the presence of the head-on collisions. The diffusive aperture for $xx$ crossing is $0.5\sigma$ higher without head-on collisions and about $3\sigma$ larger with head-on collisions. For this working point, the $yy$ crossing produces no diffusive aperture, neither with nor without head-on collisions.

![Figure 6](image)

Figure 6: Simulated diffusion rate for LHC at top energy as a function of start amplitude for $xx$, $xy$, and $yy$ crossing with long-range collisions and with the combined effect of long-range and head-on (‘short-range’) collisions for identical 0-amplitude horizontal tunes 64.30268 and a vertical tune which varies from 59.27268 to 59.31268 in steps of 0.005. The start amplitudes $A$ were chosen equal in $x$ and $y$ ($A = \bar{x} = \bar{y}$). Only the two primary IPs 1 and 5 are considered.

The dependence on the working point is illustrated in Fig. 7, which combines a large number of similar simulation results. The only difference between pictures is the vertical tune, which varies from 59.27268 to 59.31268 in steps of 0.005. The figure illustrates that the diffusive aperture is always smallest for the alternating crossing. The sole exception is the case of the coupling resonance (on the left, second row from the bottom), where the diffusion cannot be accurately determined.

![Figure 7](image)

Figure 7: Simulated diffusion rate for LHC at top energy as a function of start amplitude for $xx$, $xy$, and $yy$ crossing with long-range collisions and with the combined effect of long-range and head-on (‘short-range’) collisions for identical 0-amplitude horizontal tunes 64.30268 and a vertical tune which varies from 59.27268 (top left) to 59.31268 (bottom) in steps of 0.005. The start amplitudes $A$ were chosen equal in $x$ and $y$ ($A = \bar{x} = \bar{y}$). Only the two primary IPs 1 and 5 are considered.

Similar simulation scans were performed for a reduced crossing angle of 200 μrad (instead of the nominal 285 μrad) as well as for the nominal crossing angle and a horizontal integer tune of 63 (instead of 64). The diffusive apertures found for all these cases are displayed in Figs. 8–11 as a function of the vertical tune. We notice that for
crossing the diffusive aperture is continually low, independent of the working point. For equal-plane crossing it is almost always higher, and strongly tune dependent. At some working points, the diffusive aperture at the nominal crossing angle exceeds 14σ compared with 6σ for the alternating-crossing case. Equal-plane crossing may allow reducing the crossing to 200 μrad, while still achieving a diffusive aperture that is higher than for the nominal 285 μrad crossing angle with alternating crossing. The simulation result varies with the integer tune, since a change in integer tune changes the phase advance between the IPs, here assumed to be symmetrically spaced.

Further insight in the dynamics is provided by a frequency-map analysis à la Laskar [30, 31, 32], which was performed for the different crossing schemes. Figures 12–14 present the tune footprints for the nominal LHC working point. The shape of the footprints substantially differs for the three crossing schemes, reflecting differences in the detuning with amplitude. The horizontal-horizontal crossing schemes develops a tail in the lower right direction, overlapping the linear coupling resonance, on which some particles are trapped. The vertical-vertical crossing shows a tail towards the upper left direction, and particles trapped here on two, 11th and 13th order, resonances. The alternating crossing exhibits tails in either directions, and, consequently, particles trapped on all these resonances.

In order to illustrate the sensitivity of these features to the working point, Fig. 15–17 show the equivalent tune footprints for a vertical tune about 0.1 units lower, which corresponds to a working point on the coupling resonance. The main characteristics of the tune footprints remain unchanged.

In the simulations presented so far, the start amplitudes were always chosen at 45° in the xy plane (i.e., \( \hat{x} = \hat{y} \)), and the IPs were taken to be symmetrically positioned around the ring. We can also scan the angle \( \alpha \equiv \arctan(\hat{x}/\hat{y}) \) of the start amplitudes \( \hat{x} \) and \( \hat{y} \) in the xy plane, and, in addition, we introduce a phase-advance difference (e.g., 0.02 in units of 2π) between the two IPs in both planes.
Figure 11: Simulated LHC diffusive apertures for a horizontal tune of 63.30268 and a nominal crossing angle, as a function of the fractional vertical tune. The six curves correspond to the three crossing schemes (blue: \(xy\) crossing, red: \(xx\) crossing, green: \(yy\) crossing) with (solid) and without head-on collision (dashes). The initial starting amplitudes \(A\) were chosen equal in \(xy\) and \(xx\) and \(yy\) only.

The diffusive apertures was estimated from the amplitude for which the emittance diffusion normalized to the design emittance is \(10^{-4}\) per turn.

This case, diffusion rates as a function of starting amplitude \(r \equiv \sqrt{x^2 + y^2}\) along different angles \(\alpha\) are shown in Fig. 18. Figure 19 presents the simulated diffusive aperture as a function of the launch angle in the \(xy\) plane for the three crossing schemes. It can be seen that strongly enhanced stability is expected in the horizontal plane for \(yy\) crossing and in the vertical plane for \(xx\) crossing. Indeed, the stability limits of 20\(\sigma\) indicated below 20\(\sigma\) for \(xx\) crossing and above 70\(\sigma\) for \(yy\) crossing only signify lower bounds (maximum range of the amplitude scan). This asymmetry was expected from contour lines for the determinant of the Jacobian matrix \(\partial Q_{xy}/\partial Q_{xx,yy}\) [27, 30] and it reproduces independent simulations by W. Herr [33].

Tentatively summarizing the results of the crossing-scheme simulations studies, we found that in most cases the simulated diffusive aperture for equal plane crossing is comparable to, or larger than, that for alternating-plane crossing. In the non-crossing plane, the stability is substantially enhanced. Possible explanations for the higher stability of the equal-plane crossing are (1) a different shape of...

\footnote{In all previous plots, with \(\tilde{y}\), the amplitude was defined as \(A = \tilde{x} = \tilde{y}\), so that for \(x = y\) its value differed by a factor of 1/\(\sqrt{2}\) from the corresponding value of \(r\).}
the detuning with amplitude (‘folding’), since $x y$ crossing cancels the dodecapole and 20-pole terms in addition to the linear tune shift; and (2) an increase in the number of excited resonances for the alternating crossing. Also, particles launched in the non-crossing plane simply stay further away from the opposing beam. Nevertheless, the alternating crossing scheme ensures that the tunes of the PACMAN bunches are similar to the tunes of the nominal bunches, which is sometimes considered to be a decisive operational advantage.

2.3 Experiments at the SPS

In 2004, an experimental comparison of different crossing schemes and a benchmarking of the pertinent computer simulations became possible at the CERN SPS, with the installation of a second generation prototype beam-beam compensator [34, 35], abbreviated with the equipment code ‘BBLR’. Figure 20 shows a unit of this model, which is equipped with three copper wires, mounted vertically, horizontally and at 45 degrees, respectively. Together with the two units of the first-generation BBLR, which accommodate only a vertical wire, this device has allowed an experimental ‘simulation’ of various crossing schemes, by powering different wires with current and adjusting the transverse beam position. The first BBLR alone was used to ‘simulate’ the effect of long-range beam-beam collisions for a single-plane crossing [37].

Similarly, if only one of the two BBLRs is excited, its effect corresponds to that of the uncompensated LHC long-range collisions in one (now variable) plane of crossing. Figure 21 shows the result of such a measurement, where the beam lifetime was recorded as a function of the beam-wire separation. The dependence is perfectly fitted by a 5th order power law of the form

$$\tau \approx 5 \text{ms} \left( \frac{d}{\alpha} \right)^5. \quad (1)$$

A naive extrapolation to the typical LHC separation of $9.5\sigma$ yields a beam lifetime of about 6 minutes. This may be a pessimistic estimate, as tune ripple, intrabeam scattering, and other noise sources will be much smaller in the LHC than in the SPS. Unfortunately, since all beam-wire separation scans in the SPS were conducted for a constant beam momentum of 26 GeV/c, no experimental information on the variation of (1) with beam energy exists.
Figure 18: Diffusion rate as a function of start amplitude \( \tilde{r} = (\tilde{x}^2 + \tilde{y}^2)^{1/2} \) at \( \alpha = \arctan(\tilde{y}/\tilde{x}) \) equal to 25°, 45°, and 65°. The 0-amplitude tunes are \( Q_x = 64.30260 \) and \( Q_y = 59.26208 \). The phase advance difference between the two half rings is taken to be 0.02 (2\( \pi \)).

Note that Fig. 21 refers to a scaled experiment, in which the wire current \( I_{\text{wire}} \), normalized emittance \( \gamma \epsilon \), and beamwire separation \( d_y \) were reduced compared with the nominal set up (\( I_{\text{wire}} = 267 \, \text{A}, \, \gamma \epsilon = 3.75 \, \mu\text{m} \), \( d_y = 21.42 \, \text{mm} \)), obeying the following scaling relations:

\[
I_{\text{wire}} = 267 \frac{\gamma \epsilon}{3.75 \, \mu\text{m}} \Lambda
\]

Figure 19: Diffusive aperture, approximated by the start amplitude \( \tilde{r} = (\tilde{x}^2 + \tilde{y}^2)^{1/2} \) at which the diffusion rate in units of design emittance exceeds \( 10^{-3} \) per turn, as a function of starting angle in \( xy \) plane, \( \alpha = \arctan(\tilde{y}/\tilde{x}) \). The 0-amplitude tunes are \( Q_x = 64.3026 \) and \( Q_y = 59.26208 \). The phase advance difference between the two half rings is taken to be 0.02 (2\( \pi \)).

Figure 20: Photo of the second-generation ‘BBLR’ in the CERN SPS (G. Burtin, J. Camas, J.-P. Koutchouk, et al.)

\[
d_y = 21.42 \sqrt{\frac{\gamma \epsilon}{3.75 \, \mu\text{m}}} \, \text{mm}.
\]

The advantages of this scaling are an increase in the physical aperture compared with the dynamic aperture, and an improved beam lifetime for the ‘bare’ machine.

In a similar experiment at the Tevatron, using the electron lens (TEL) as a ‘wire’, the beam lifetime was measured as a function of the beam-TEL distance [38]. The Tevatron lifetime increased as the third power of the beam-TEL distance (see Fig. 22). The different power law might be related to the different working points in tune diagram. The antiproton tunes in the Tevatron were about \( Q_x \approx 0.58 \) and \( Q_y \approx 0.59 \), while the ‘LHC-like tunes’ for the SPS experiments were \( Q_x \approx 0.32 \) and \( Q_y \approx 0.29 \). Another difference between the two experiments is that at the SPS the tunes were corrected and held constant for each beam-wire distance, while at the Tevatron the tunes changed (slightly).
Figure 21: Fitted SPS beam lifetime as a function of the beam-wire distance, in an experiment conducted during the HHH-2004 workshop on 09/11/2004, with a single-wire excitation of 190 A, tunes $Q_x = 0.3208$, $Q_y = 0.2914$, emittance $\gamma_{x} \approx \gamma_{y} \approx 1.4 \mu m$, and an rms beam size of about 1.56 mm. The beam was vertically separated from the wire, at the same horizontal position.

Figure 22: Antiproton loss rate as a function of the TEL distance, exhibiting an inverse cubic dependence (V. Shiltsev). The TEL current was about 600 mA, the electron-beam transverse rms size 0.66 mm, and the TEL length 2 m. The antiproton beam energy is 980 GeV. A proton beam was also present during the measurement, separated about 7 mm from the antiproton beam [38].

during the scan. However, a plausible explanation of the observed low powers is missing. Higher order resonances which are supposed to dominate at the working points of the two experiments should have given rise to powers of higher order.

The experiments described above employed only one wire. A first experiment on the crossing planes was conducted in the CERN SPS on the 26th of August 2004, with the nominal emittance. Three configurations were implemented as illustrated in Fig. 23. Due to constraints by the physical aperture and the different distances of the horizontal and vertical wire from the center of the chamber (about 55 mm and 20 m, respectively), a pure alternating crossing could not be realized. Instead a mixed scheme was chosen (configuration a), modeling horizontal crossing at one wire and 45-degree crossing at the other, by exciting both wires at the same current. Equal-plane crossings were modeled by exciting only one of the two wires at twice its original strength (configurations b and c). For completeness, and to observe a larger effect on the beam lifetime, the first configuration was also tested at twice the strength, which corresponds to a two times higher intensity of the opposing LHC beam.

The dynamic aperture for the three (four) configurations were simulated by the WSDIFF code as a function of the vertical tune for a constant horizontal tune of $Q_x = 0.31$. The result of the simulated tune scan is shown in Fig. 24. It can be seen that the dynamic aperture of the horizontal-horizontal crossing is about 10% larger than for the mixed (or “quasi-XY”) and the 45-degree (or “quasi-YY”) configurations. The reduction in dynamic aperture for a doubling in the wire current is at the most 1/σ.

In the experiment, the beam lifetime was measured for different vertical tunes, as shown in Fig. 25. For most tune values, also in the experiment the horizontal-horizontal crossing exhibited the best beam lifetime, the pure 45-degree crossing the second best, and the mixed crossing the worst. However, the situation changed towards the two ends of the scan range, near the 7th and 3rd integer resonance, respectively. Here, the pure 45-degree crossing scheme was the most robust, whereas for all other configurations the lifetime strongly decreases. It may be that lattice nonlinearities are responsible for this drop in lifetime, and we could conjecture that the 45-degree cross-
Figure 24: Dynamic aperture simulated by the WSDIFF code as a function of the vertical tune for the three SPS wire configurations shown in Fig. 23. The horizontal tune is held constant at $Q_{\chi} = 0.31$.

Figure 25: Beam lifetime measured as a function of the vertical tune for the three SPS wire configurations of Fig. 23. The horizontal tune was held constant at $Q_{\chi}$.

Since in the first SPS experiment, described above, the bare beam lifetime was comparable to the lifetime with one or two wires excited, we performed a second experiment for which we reduced the beam-wire distances and used smaller emittances. We applied the scaling of Eqs. (2). One of the wires (BBLR2) had been rotated prior to the experiment, in order to allow for shorter transverse distances. The three configurations of Fig. 26 could then be realized. Again, it was not possible to implement a pure horizontal-vertical crossing. However, instead a pure $45^\circ$-$135^\circ$ inclined hybrid crossing [26] was realized and the performance could be compared with that of a pure vertical-vertical or pure $45^\circ$-$45^\circ$ crossing. Because the distances were smaller than in the preceding experiment, the total wire current was reduced to 190 A (to be compared with 240 A used in the previous study, which had been close to the maximum current allowed by wire heating.) Since the bump amplitude and beam emittance were smaller, the relative size of the physical aperture was increased in this experiment, which improved the natural beam lifetime and was expected to render the effect of the wires more easily discernible.

Figure 26: Approximations of different crossing schemes in the second SPS experiment. The second BBLR was rotated by 180 degrees. The beam was bumped horizontally and vertically. The emittance was reduced to about $\gamma e = 1.4 \, \mu m$, and either one or two wires were excited at different strengths. Configuration a) models $45^\circ$-$135^\circ$ inclined hybrid collision [26], b) a double $45^\circ$ hybrid crossing and c) a pure vertical-vertical crossing.

Figure 27 displays the simulated dynamic aperture for the three configurations of the second experiment. The pure $45^\circ$-$45^\circ$ crossing has the smallest dynamic aperture. At vertical tunes of 0.29 or lower the vertical-vertical crossing is best, while at higher tunes the inclined-hybrid scheme yields the largest dynamic aperture. For completeness, the simulation results for a pure horizontal-vertical crossing are also indicated. On average its dynamic aperture is about the same as that for the inclined-hybrid scheme, but the two are not identical.

The experimental beam lifetimes as a function of vertical tune are presented in Fig. 28. The measured beam lifetime is lowest for the $45^\circ$-$45^\circ$ crossing, consistent with the simulations. Also, that the inclined hybrid crossing is best for tunes above 0.3 and the pure vertical-vertical crossing for lower tunes, agrees well with the simulation in Fig. 27.

3 ELECTRON LENSES

The basic idea of the Tevatron electron lens (TEL) [39, 40, 41, 42] is “to compensate (on average) space charge forces of positively charged protons acting on antiprotons in the Tevatron by interaction with a negative
Figure 27: Dynamic aperture simulated by the WSDIFF code as a function of the vertical tune for the three SPS wire configurations shown in Fig. 26. The horizontal tune is held constant at $Q_x = 0.32$.

Figure 28: Beam lifetime measured as a function of the vertical tune for the three SPS wire configurations of Fig. 26. The horizontal tune was held constant at $Q_x = 0.32$.

Figure 29: Schematic of antiproton beam-beam tune shift compensation at the Tevatron using an electron lens [42].

Figure 30: Drawing of the Tevatron Electron Lens (TEL) with some of its components. Not shown are the high-voltage modulator, high-voltage and high-current power supplies, cryogenics, solenoid quench protection, vacuum pumps, control system, diagnostics, and cables [42].

Figure 31: Photo of the Tevatron Electron Lens (TEL).

charge of a low-energy high-current electron beam” [42]. A schematic of the TEL principle is shown in Fig. 29. Figure 30 presents a schematic of the device itself and Fig. 31 a photo. The solenoid field is 3.5 kG, limited by its steering effect on the electron beam. The maximum kinetic energy of the electrons is 15 keV, and the pulsed electron current can be raised up to 6 A, with 0.6–1.0 A being a typical value. The effective length of the TEL is 2 m and the rms electron beam radius about 0.66 mm. The first TEL is in operation since 2001. A second TEL was fabricated in collaboration with IHEP Protvino. This TEL will be installed during 2006, at a location with a different ratio of horizontal and vertical beta function, so that the two TELs will provide an independent control of the two transverse tunes.

As a historical aside, we mention that E. Tsyganov and co-workers had in the early 1990s proposed a similar set up, consisting of a mixture of 10-keV and 1-MeV electron beams in a 2-T solenoid, for compensating the tune spread arising from the head-on beam-beam interaction in the SSC and the LHC [10, 11].

The FNAL TEL is pulsed and can be adjusted individually for each antiproton bunch. Thereby, the bunch-to-bunch tune variation can be minimized. By properly tailoring the transverse profile of the electron beam so as to ‘match’ that of the proton beam, also the intra-bunch tune spread may be reduced. For the Tevatron, the intra-bunch tune spread and inter-bunch tune variation are of
roughly comparable magnitude, as is illustrated by Fig. 32. The electron lens may be used not only in collision, but also at injection, where long-range beam-beam effects have proven to be an issue as well [17], e.g., for equalizing the bunch tunes. The electron-beam current is variable so that the TEL can track a changing proton beam current.

The tune shift induced by the TEL is described by

$$\Delta Q_{x,y} = \pm (-1) \frac{\beta_x \gamma_p}{2\pi} \frac{1}{\beta_e} \frac{I_e I_{\text{TOT}}}{\epsilon_{\text{TOT}}} \gamma_p \gamma_p,$$

(3)

where $I_e$ is the electron current, $\alpha_e$ the total electron-beam radius (assuming a uniform profile), $I_x$ the length of the beam-electron interaction, $\beta_e$ the electron velocity divided by the speed of light, $\gamma_p$ the relativistic Lorentz factor of the antiproton (or proton) beam, and the two signs refer to the situations that the electron beam moves in a direction opposite to the beam to be affected (antiprotons) or in the same direction as this beam (protons). Tune shifts of 0.01 at 150 GeV are readily obtained.

The beam-beam interaction also drives resonances. Compensating resonance driving terms by the FNAL TEL would require a great optics flexibility, e.g., allowing for the adjustment of the phase advances from the proton-collision points to the TEL and of the beta functions at the TEL. Therefore, a resonance compensation is not within the scope of present TEL experiments.

Two features of the electron lens have proven critical in the experiments conducted so far [43]. The first is the transverse profile of the electron beam. A sharp edge introduces strong nonlinear forces and ‘scrapes off’ any antiprotons oscillating at amplitudes larger than the electron-beam radius. This effect was pronounced for the first electron gun employed, which produced a beam with a uniform profile. The beam lifetime was much improved by a second gun delivering a beam of transversely Gaussian shape. A third gun which combines a uniform flat top and Gaussian tails is under development. The second critical factor is the alignment of the electron beam to the antiproton beam. This alignment is presently compromised by the different lengths of the antiproton bunches (2 ns rms, with a bunch spacing of about 396 ns) and the TEL electron pulse (total pulse duration between 0.1 and 1.0 $\mu$s). The reading of the beam-position monitors (BPMs) was found to shift by about 1.5 mm as a function of pulse length [43]. A new type of BPM will be installed to mitigate this problem.

At the Tevatron a fast emittance growth frequently occurs at the start of high-energy collisions, especially for the weaker antiproton beam. Due to the 3-fold symmetry of the Tevatron bunch filling scheme, the emittance growth is nearly the same for the corresponding bunches in each of three 12-bunch trains. Therefore, the emittance growth for only the first antiproton train is shown in Fig. 33. The blow-up rates are smaller for the bunches at the start and end of a train (namely the ‘PACMAN bunches’). This phenomenon of rapid emittance growth at the center of a bunch train is known by the name ‘scallops’, due to its appearance on the control-room screen display. The scallops effect is caused by the beam-beam interaction. Its occurrence and severity depend on the antiproton tune, and, specifically, on the proximity to 5th, 7th and 12th order resonances. Tune changes by $2 \times 10^{-5}$ can alter the behavior [42].

The year 2003 saw the first successful demonstration of TEL beam-beam compensation. In an experiment, the TEL slowed down the emittance growth of a specific antiproton bunch, when fast emittance growth was present. The successful compensation attempt is illustrated in Fig. 34. The TEL acted on bunch #33, whose emittance growth of 1.0 $\pi$ mm-mrad/hr is smaller than that of the other two equivalent bunches (4.1 and 2.2 $\pi$ mm-mrad/hr, respectively). The expected TEL tune shift for the experimental conditions was 0.001–0.004 [42]. Other similar experiments had shown ambivalent or slightly negative results, indicating that the effect of the TEL is not well controlled. An explanation for this lack of control is that the precise centering of the electron lens and the antiproton beam is difficult due to the differences in the BPM readings for the electron beam and the antiproton bunches discussed above. The BPM problem is aggravated by a migration of the antiproton orbit by about 0.5 mm over 12 hr [43].

4 ELECTROMAGNETIC WIRES

An elegant compensation of the effect of the long-range beam-beam collisions in the LHC by means of current-carrying wires has been proposed by J.-P. Koutchouk [34, 36]. The field of a wire running parallel to the beam depends on the radial distance from the wire, $r$, as $1/r$, exactly as the combined electric and magnetic force exerted by the opposing beam. For the LHC, a compensation based on such a wire is facilitated by the fact that all long-range collisions on one side of an IP occur at nearly the same

Figure 32: Tune distribution for 36 antiproton bunches colliding with 36 proton bunches in the Tevatron, evidencing both the intra-bunch tune spread and the tune variation between bunches, as computed by Yu. Alexahin and presented by V. Shiltsev [42]. The goal of the ideal electron lens is to squeeze all particle tunes into a single point.
betatron phase. Therefore, a single wire placed at about the same phase allows the simultaneous compensation of a large number of long-range collisions. This compensation being local, \textit{i.e.}, close to the source of the disturbing effect, it can, in principle, correct all nonlinear (as well as the linear) effects of the long-range collisions.

To achieve the local compensation, the wire must be positioned at a location where the two beams are already separated, since it otherwise would interfere with the second beam. For the compensation to be efficient, the transverse distance between the wire from the beam in units of the rms beam size should be about the same as the typical distance of the long-range beam-beam collisions. The proposed location is upstream (or downstream) of the LHC dipoles D1 which combine (or separate) the two beams as they approach (or leave) the interaction region. At top energy, the average phase advance between the foreseen location of the compensator and the long-range collision points is 2.6°. According to simulations [45], this is sufficiently small to ensure a satisfactory compensation. For the same transverse distance, the integrated strength of the wire (excitation current $I_{\text{wire}}$ times length $L_{\text{wire}}$) must equal the product of the opposing beam bunch charge $N_{\text{be}}$ ($N_{b} \approx 1.15 \times 10^{11}$) and the number of long-range collisions encountered ($n_{\text{par}} \approx 15$ for a nominal bunch at the center of a bunch train), multiplied with the speed of light: \begin{equation} I_{\text{wire}} L_{\text{wire}} = N_{\text{be}} c n_{\text{par}}. \end{equation}

Note that the force due to the wire is purely magnetic, while the other beam produces a magnetic and an electric field. For an ultra-relativistic beam, these two forces are equal and additive. The resulting factor of two is cancelled against a factor 1/2 arising from the reduced time of interaction of the beam-beam collision, due to the opposite direction of motion of the two beams. The wire current required for nominal LHC parameters is about 82 A. It increases to about 120 A for ultimate LHC parameters (a bunch population of about $N_{b} \approx 1.67 \times 10^{11}$).

There are two limitations to the wire-compensation scheme. First, the compensation ceases to work for large betatron amplitudes when particles approach the center of the opposing beam, where the beam field no longer has the simple $1/r$ form of the wire. The maximum amplitude at which the compensation can work may roughly be estimated as the beam-beam distance minus two rms beam sizes. It decreases slightly for higher beam currents. The situation could be improved by replacing the wire with a more intricate sheet current, which may reproduce part of the deviation from the pure $1/r$ law due to the finite size of the opposing beam. However, it is well known that any finite-term expansion of the beam-beam force around the origin leads to large divergences. So, any wire-based compensation is bound to fail at amplitudes smaller than the distance to the other beam, and the gain of a more complicated scheme may therefore be marginal. A second limitation is that a dc wire would exactly compensate the long-range beam-beam effect only for nominal bunches. In the LHC these are about half the total number. All other bunches — especially those at the start and tail of a bunch train — belong to one or another class of ‘PACMAN’ bunches, which do not experience the full number of long-range collisions. To compensate the long-range beam-beam effect also for these PACMAN bunches, the wire current would need to be pulsed, \textit{e.g.}, ramped up at the start of a 2-µs bunch train over a time interval of 375 ns and down at the end of a train. Depending on the spacing between bunch trains, the ramp-down would go to zero or to an intermediate value and the ramp-up or ramp-down times change accordingly. This indicates that the time pattern of the excitation must be sufficiently flexible, so that it can be adapted to different operating conditions.

Tune calculations by MAD have shown that the tune spread of about $2 \times 10^{-3}$ for betatron amplitudes up to $6\sigma$, induced by long-range collisions at one IP, can be reduced by at least a factor of 10 using the proposed compensating wire [34].
To study the feasibility of a wire-based beam-beam compensation, an experimental programme with compensator prototypes has been conducted in the CERN SPS since 2002, when a single wire device (‘BBLR’) was installed. An SPS BBLR consists of two adjacent units, each with an effective length of 60 cm, connected in series. The wire has an outer diameter of 2.54 mm. Its copper wall is 0.5 mm thick. The wall surrounds an inner hollow channel fed with cooling water. The total physical length of a 2-unit device is 1.85 m. The first BBLR was used, during 2 years, to ‘simulate’ experimentally the effect of LHC long-range collisions and to benchmark the simulations. Figure 35 illustrates that the simulated diffusion due to the wire indeed resembles the simulated diffusion from LHC long-range beam collision in this scaled experiment. Most BBLR experiments were performed at 26 GeV/c, and one at 55 GeV/c. The higher beam energy is preferred, since it increases the relative size of the physical aperture compared with the dynamic aperture which is to be measured.

The experimentally observed linear-optics effects of the BBLR, like closed-orbit orbit distortion and tune shifts induced by the wire, agree with analytical calculations and simulations, and they can be exploited to precisely monitor the beam-wire distance. We note that a self-consistent solution for the tune and the closed-orbit shift is required as the change of one affects the other [37]. The non-linear effects of the wire manifest themselves in beam losses, in a reduction of the dynamic aperture, as evidenced by a shrinking emittance and by a degradation of the beam lifetime, in a decreased decoherence time and in a change of the detuning with amplitude. Several scans were performed of the wire current (corresponding to beam current in LHC) and of the beam-wire distance (corresponding to the LHC crossing angle). An example of a more recent distance scan was given in Fig. 21. The results of the 2002–2003 single-wire experiments in the SPS were summarized in two reports [37, 47]. They are not reproduced here.

For 2004, two new BBLR devices were built, each equipped with three wires, in order both to explore the compensation efficiency and to compare the consequences of alternative crossing schemes (see Section 2.3 above). Only one of the two new BBLRs could be used in 2004. It was mounted close to the existing first-generation BBLR separated by a phase advance of about 2.6°, which equals the expected average phase advance between compensator and long-range collision points in the LHC. Improper welding had caused a vacuum leak in the second new device, which is presently being refurbished and which will be installed on the other side of the SPS, the ‘long straight section 2’, during the long 2005 shutdown. A photo of a new 3-wire BBLR was shown in Fig. 20. Figure 36 presents a cross section of this BBLR, which also indicates that it is movable in the vertical direction over a total range of 5 mm via remote control. The vertical mobility has allowed optimizing the distance to the beam for maximum compensation and studying pertinent tolerances. The compensation effect was found to be fully lost for a movement of about 3 mm, as predicted by simulations [48].

![Figure 36: Cross section of 3-wire BBLR device installed in SPS LSS5 since July 2004 (Courtesy G. Burtin).](image-url)
Figure 37 shows the range of a vertical tune scan around the LHC tune on 30/07, and Fig. 38 presents the associated beam lifetimes for the three situations of the bare machine (no BBLR excitation), with only one BBLR excited, and with the compensating BBLR also turned on. It can be seen that the single BBLR reduced the observed beam lifetime from about 250 s to 50 s. The compensation restored the original lifetime nearly perfectly, except for at tunes below 0.285. This coincides with the tune region where the 7th and, below 0.275, also the 4th order resonances are found. On the other hand, the imperfect compensation for low tunes could also be an artifact of drifting machine conditions during the experiment. A future repetition of the experiment may clarify this point. In all cases the beam lifetime drops as the tune approaches the 3rd integer resonances, at 0.32–0.33. This lifetime reduction must be driven by the natural nonlinearities of the SPS, since it is only weakly affected by the BBLR.

Figure 39 shows a result from the second ‘scaled’ experiment on 02/09/04. Here the tune was held constant. The beam lifetime is shown, evolving in time as the conditions of the experiment were changed. Note that the lifetimes were a factor 20 longer than for the previous experiment, with larger emittances. The compensation not only restored, but even improved the lifetime compared to that of the bare machine. Perhaps a small additional nonlinearity from the BBLRs has a stabilizing effect, or, else, also this overcompensation may be an artifact from drifts in the PS or SPS beam conditions. Such drifts or step changes are sometimes noticed, e.g., as a sudden variation in beam emittance, but their origin is not known (often attributed to the PS which provides the injected beam and where a delicate rf gymnastics is performed). They cannot be well controlled or even fully characterized during the experiment.

Figure 37: Tune path traced during the first SPS wire-compensation experiment on 30/07/04., and resonance lines through 11th order.

Figure 38: Beam lifetimes measured as a function of the vertical tune without BBLR (blue diamonds), with one BBLR excited at 240 A (green triangles) and with both BBLRs excited so that they compensate each other (red squares), on 30/07/04.

Figure 39: Beam lifetime as a function of time with compensation (left), without compensation (center) and for the bare machine (right), measured on 02/09/04.

5 SUPER-BUNCHES

The concept of a ‘super-bunch hadron collider’ was proposed by K. Takayama and co-workers in 2002 [26]. A schematic showing several long superbunches colliding in alternating planes at two interaction points is given in Fig. 40. An underlying feature of the super-bunch scheme is an intrinsic cancellation between parts of the long-range and head-on beam-beam tune shifts experienced at the two interaction points, which exploits that the sign of the long-range tune shift in the plane of the crossing is opposite to the sign of the head-on tune shift.

Studies of luminosity optimization for an LHC upgrade [49, 50, 51, 52] pointed into a similar direction. Namely, as shown in Refs. [51, 52], for (1) a large Piwinski angle

\[
\phi_{T5} \equiv \frac{\theta_{\epsilon} \sigma_d}{2\sigma_{\epsilon}} \ll 1 , \tag{5}
\]

...
(2) a crossing angle much smaller than $\pi/2$, but larger than the IP rms beam divergence, i.e., $1 \ll \theta_c \ll \sqrt{\sigma_z/\beta^*}$, and
(3) an rms bunch length large compared with the transverse IP beam size ($\sigma_z \ll \sigma^*$), the luminosity expression for long round Gaussian bunches colliding at two interaction points in alternating planes under a full crossing angle $\theta_c$ can be simplified to
\[
L_{\text{Gauss}} \approx \frac{1}{2} \frac{f_{\text{coll}} \gamma^2}{r_p \beta^*} \Delta Q_{\text{tot}} N_b ,
\]  
while that for round bunches with uniform 'flat' longitudinal profile transforms to
\[
L_{\text{flat}} \approx \frac{1}{\sqrt{2}} \frac{f_{\text{coll}} \gamma^2}{r_p \beta^*} \Delta Q_{\text{tot}} N_b .
\]
Comparing (6) and (7) reveals that flat bunches can deliver higher luminosity. Namely, for the same total beam-beam tune shift $\Delta Q_{\text{tot}}$ (sum over the tune shift encountered at the two interaction points) and the same bunch charge $N_b$ (in units of the electron charge), the luminosity is a factor of $\sqrt{2}$ higher for the flat longitudinal profile. In the above equations $f_{\text{coll}}$ denotes the bunch collision frequency, $\beta^*$ the IP beta function, taken to be the same in both transverse planes, and $\gamma$ the particle energy divided by the rest mass.

Equation (7) applies not only to super-bunches, but also to conventional bunches with a flat profile, as long as their corresponding Piwinski angle is large. Figure 41 shows a schematic illustration of the luminosity gain achievable by increasing crossing angle or bunch length and by flattening the profile. Evaluating the more precise luminosity expressions of [49] and augmenting them by bunch-end effects [53], H. Damerau showed that in case of the LHC the full factor $\sqrt{2}$ gain in luminosity from the shape of the beam profile can be realized for Piwinski angles $\phi_{\text{Pw}}$ larger than about 10 [53].

6 CRAB CAVITIES

The crab cavity was first proposed by R. Palmer to increase the luminosity at linear colliders [54]. Shortly thereafter, the concept was extended to storage-ring colliders [55]. The first installation of crab cavities in an operating collider is foreseen for the end of 2005 at the KEK B factory [56].

The basic idea of crab crossing is illustrated in Fig. 42. The differential deflections received in dipole-mode rf cavities, located upstream of the IP and the low-beta quadrupoles, align the bunches at the collision point, so that the geometric luminosity is the same as for head-on collisions. The head-tail motion of the bunch is removed by a second equivalent crab cavity on the outgoing side of the IP. The rf phases of the incoming and outgoing crab cavities must be held stable with respect to each other, since any change in their phase difference leads to a bunch centroid position at the IP.

The effective head-on collision provided by the crab

![Figure 40: Schematic of Super Bunches in a High-Luminosity Collider [26].](image)

![Figure 41: Relative increase in LHC luminosity vs. relative increase in rms length (or crossing angle), for Gaussian or hollow bunches, maintaining a constant beam-beam tune shift with alternating crossing in two interaction points. The luminosity is taken to be $\sqrt{2}$ higher for the flat beam profile, as predicted by the asymptotic equations (6) and (7). The bunch population increases in proportion to the luminosity gain. The number of bunches is unchanged.](image)

![Figure 42: Schematic of crab crossing principle.](image)
crossing avoids the excitation of synchro-betatron resonances and may allow for higher beam-beam tune shift [57]. At the same time, the centre of mass of the bunches follow a trajectory with crossing angle, so that the beams can be easily separated without suffering from long-range collisions.

Hence the crab cavity combines all advantages of both the head-on and crossing-angle schemes: (1) high luminosity, (2) high beam-beam tune shift limit, (3) weak long-range collisions, (4) large crossing angle, and (5) simplified IR design and relaxed parameters of the low-beta quadrupoles.

A challenge may be the crab cavity rf system, which should provide sufficient crab voltage, but not blow up the beam emittance. The required crab rf deflecting voltage is

\[
V_\perp = \frac{c E_b \tan \theta_\perp/2}{\epsilon \omega \sqrt{\beta_x \beta_y^{\text{crab}}}},
\]

where \(E_b\) is the beam energy, \(\omega\) the angular rf frequency, and \(\beta_x\) the beta function at the crab cavity.

If the crab cavity rf phase fluctuates with respect to the crab cavity of the opposing beam at the other side of the IP, a centroid displacement \(\Delta x\) between the two beams at the collision point results. At the LHC the turn-by-turn jitter in the displacement between the two beams, \(\Delta x\), should be smaller than 12 nm (about 0.1% of the IP beam size) in order to ensure an emittance growth of less than 10% per hour. This limit \(\Delta x_{\text{max}}\) translates into a limit for the acceptable crab rf phase jitter of [58]

\[
\Delta \phi \leq \frac{\Delta x_{\text{max}} 2\pi}{\lambda \epsilon \theta_\perp}.
\]

Note that this is not an absolute stability requirement, but a constraint on the relative jitter for the crab cavities of the two beams on either side of the IP.

Table 3 compares the crab-cavity parameters required for an LHC upgrade with those of the KEKB factory. The table assumes a crossing angle of 8 mrad, which is a factor 4 larger than the minimum angle for which the beams may still be fed into two separate low-beta quadrupoles (see Refs. [59, 60]). The LHC requires about 100 times more deflecting voltage than KEKB, primarily due to the increased beam energy. The turn-by-turn phase-jitter tolerances are of the order milli-degrees. Note also that the rf frequency of 1.3 GHz would be too high for the present nominal LHC bunch length.

### 7 CONCLUSIONS

We have discussed a number of approaches for cancelling or compensating undesired aspects of the beam-beam interaction and, thereby, optimizing the luminosity of a hadron collider.

Alternating the plane of crossing at two (or more) interaction points cancels the linear tune shift induced by the long-range collisions on either side of an IP, but seems to provide little gain in dynamic aperture. In several cases, the latter even appears to be reduced by this type of compensation. Nevertheless it is the baseline scheme for the LHC, since it also equalizes the tunes of so-called PAC-MAN bunches, at the head or tail of a bunch train, with that of the nominal bunches.

Programmes of active beam-beam compensation are in progress for the Tevatron and the LHC.

The Tevatron Electron Lens is a highly promising approach, with a successful improvement of the beam lifetime already demonstrated in machine studies. Presently the TEL still suffers from a lack of control of critical parameters, such as the relative alignment of electrons and antiprotons, which so far prevents TEL use in routine operation.

Wire compensation of the long-range collisions at the LHC will allow for smaller crossing angles and higher bunch charges. An experiment at the SPS has established the feasibility of a dc compensation. A pulsed-wire operation would be desirable for a selective correction of the PACMAN bunches.

Flat long bunches or extreme super-bunches may yield higher luminosity for the same beam-beam tune shift (though these options are disliked by the particle-physics experiments, since they imply an increased event pile up in the detectors).

Crab cavities are a promising alternative option enabling the use of large crossing angles.

### 8 ACKNOWLEDGEMENTS

I would like to thank Jean-Pierre Koutchouk, Werner Herr, Kazuhiro Ohmi, Yannis Papaphilippou, Francesco Ruggiero, Tanaji Sen, and Vladimir Shiltsev, for many stimulating discussions and much helpful advice on the topic of this presentation. The experimental studies of wire-based beam-beam compensation in the CERN SPS were made possible by some of them, in particular by Jean-Pierre Koutchouk, as well as numerous further colleagues, including Gerard Burtin, Jacques Camas, Jean-Jacques Gras, Roger Perret, Federico Roncarolo, Marcel Royer, and Jorg Wenninger.
9 REFERENCES

[35] The BBLR installation was planned and performed by G. Burtin, J. Camas, J.-P. Koutchouk, with assistance from many colleagues. The beam experiments were conducted by J.-P. Koutchouk, J. Wenninger, and F. Zimmermann, with help from F. Roncarolo, T. Sen, V. Shiltsev, and Y. Papaphilippou.
GENERA TION AND BENEFITS
OF LONG SUPER-BUNCHES IN THE LHC

H. Damerau ∗, R. Garoby, CERN, Geneva, Switzerland

Abstract

To maximize the luminosity of the Large Hadron Collider (LHC) bunches with a uniform longitudinal line density are considered for a future upgrade. In the first part three different options will be presented: short rectangular bunches held by multiple RF harmonics using multiples of 400 MHz, long and flat bunches held by multiples of 40 MHz and the so-called super-bunches which are confined by barrier buckets. The potential luminosity improvement of these options with respect to the beam-beam limit will be analyzed. The second part will be dedicated to the generation of the bunches with uniform line density by appropriate RF manipulations. It will be shown that especially the scheme to create long and flat bunches at multiples of 40 MHz relies on well proven RF techniques and can be performed in the LHC with additional RF hardware of reasonable size and performance.

INTRODUCTION

In the nominal scheme [1], both rings of the LHC will be operated with 2808 bunches each, kept in the buckets of an RF system working at 400.8 MHz (35640th harmonic, \(h\) of the revolution frequency). To increase the luminosity by an order of magnitude beyond \(10^{34} \text{cm}^{-2} \text{s}^{-1}\) it has already been suggested to operate the machine with super-bunches [2], sausage-like bunches with a uniform longitudinal line density held by a sophisticated wide-band RF system based on induction cell cavities [3].

At a constant strong beam-beam tune shift, the operation with rectangular bunches of equal longitudinal peak density produces up to \(\sqrt{2}\) more luminosity than the collision of Gaussian bunches [4]. In the first part the analytical expressions to calculate the luminosity and beam-beam tune shift for rectangular bunches of arbitrary length are derived. These formulae are applied to three rectangular bunch options in the second part: Short rectangular bunches held by multiples of 400.8 MHz, flat bunches of intermediate length held by multiples of 40.08 MHz and super-bunches held by a wide-band RF system generating barrier buckets. Especially the first option cannot be treated by classical coasting beam formulae [5] as the length of the bunches is within their length of direct interaction.

An approach to generate flat bunches of intermediate length by means of RF manipulations is presented in the last part. This scheme allows the combination of bunches of 16 or 32 nearly nominal bunches with the nominal bunch spacing of 25 ns to a high intensity bunch held by a multiple harmonic RF system to approximate a rectangular particle distribution.

LUMINOSITY FOR RECTANGULAR BUNCHES

The Lorentz invariant differential interaction rate is defined by [6]

\[
d\frac{d^2N}{dt\,dV} = \varsigma n_1 n_2 \sqrt{(\bar{v}_1 - \bar{v}_2)^2 - \left(\frac{\bar{v}_1 \times \bar{v}_2}{c^2}\right)^2} = \varsigma \frac{d\mathcal{L}}{dV},
\]

where \(n_1 = n_1(x_1, y_1, z_1, t)\) and \(n_2 = n_2(x_2, y_2, z_2, t)\) denote the particle densities of two bunches with the velocities \(\bar{v}_1\) and \(\bar{v}_2\). The cross section is given by \(\varsigma\). It can be simplified in the case of ultra-relativistic beams crossing under angles \(\theta \ll 1\) so that the general expression for the luminosity \(\mathcal{L}\) reduces to

\[
\mathcal{L} = 2c f_0 \cos^2(\theta/2) \int n_1 n_2 \, dV \, dt,
\]

which actually represents the time and space convolution integral of the two bunch densities. For \(k\) bunches per beam the revolution frequency \(f_0\) has to be replaced by the collision frequency \(f_{\text{coll}} = k f_0\).

The density distribution function of a transversely Gaussian beam with a rectangular longitudinal line density can be written as

\[
n(x, y, z, t) = \frac{N}{l_b} \frac{1}{2\pi \sigma_x \sigma_y} \exp \left[-\frac{1}{2} \left(\frac{x^2}{\sigma_x^2} + \frac{y^2}{\sigma_y^2}\right)\right] \zeta(z),
\]

with the longitudinal distribution function being defined by a combination of two unit step functions

\[
\zeta(z) = \begin{cases} 
1 & -l_b/2 \leq z \leq l_b/2 \\
0 & \text{elsewhere}
\end{cases}.
\]

The bunch intensity is described by the number of particles \(N\).

For a convenient calculation of the luminosity integral Eq. (2) a coordinate system symmetrically tilted by \(\theta/2\) with respect to both beams is chosen (see Fig. 1) so that

![Figure 1: Choice of the coordinate system for luminosity calculation.](image)

\[ \text{beam 1} \quad \text{beam 2} \]

\[ \theta/2 \quad \theta/2 \quad \text{interaction point} \]

\[ \phi \]

* On leave of absence from Gesellschaft für Schwerionenforschung mbH (GSI), Darmstadt, Germany
the transformations become
\[ x_1 = (x \cos(\phi/2) - z \sin(\phi/2), y, x \sin(\phi/2) + z \cos(\phi/2)) \]
and
\[ x_2 = (-x \cos(\phi/2) - z \sin(\phi/2), y, x \sin(\phi/2) - z \cos(\phi/2)) \]

As this will be the most relevant case in the LHC, only round beams \( \sigma_x = \sigma_y = \sigma \) are considered below. Furthermore, both crossing bumps are assumed to be equal so that \( N_1 = N_2 = N \), \( \sigma_1 = \sigma_2 = \sigma \) and \( l_1 = l_2 = l_b \).

After execution of the integrations over \( x \), \( y \) and \( t \) the luminosity integral becomes
\[
\mathcal{L} = \frac{f_0 N^2}{2 \pi l_b} \cos(\theta/2) \int_{-l_{\text{det}}/2}^{l_{\text{det}}/2} \frac{1}{\sigma(z)^2} \exp \left( -z^2 \sin^2(\theta/2)/\sigma(z)^2 \right) \left( 1 - \frac{2|z| \cos(\theta/2)}{l_b} \right) dz ,
\]
for \( 2|z| \cos(\phi/2) \leq l_b \). The integrand vanishes elsewhere and the detector covers \( \pm l_{\text{det}}/2 \). It is important to point out that the second factor which arises from the integration over \( t \) according to
\[
\int \zeta(z_1 - ct) \zeta(z_2 - ct) \, dt = \begin{cases} \frac{l_b}{c} & \text{if } |z_1 - z_2| \leq l_b \\ 0 & \text{elsewhere} \end{cases},
\]
which contains the correction for bunch lengths in the order of the length of the luminous region or below and becomes unity for coasting beams. The beta function variation along the interaction region results in an RMS beam radius is given by \( \sigma(z) = \sigma^*(1 + z^2/b^2)^{1/2} \), where \( \beta^* \) and \( \sigma^* \) are the beta function and the beam radius at the interaction center. Assuming small crossing angles so that \( \sin \theta = \theta \) and \( \cos \theta = 1 \) as well as \( \beta^* \theta \gtrsim 10 \sigma^* \), the luminosity is approximated to
\[
\mathcal{L} = \frac{f_0 N^2}{2 \pi l_b} \frac{\beta^*}{\sigma^*} \int_{-l_{\text{det}}/(2\beta^* \sigma^*)}^{l_{\text{det}}/(2\beta^* \sigma^*)} \exp \left( -\frac{\beta^2 \theta^2}{4 \sigma^* u^2} u^2 \right) du = \frac{8}{l_b \beta^2} \left[ 1 - \exp \left( -\frac{\beta^2 l_{\text{det}}^2}{16 \sigma^*} \right) \right].
\]
The first term represents again the coasting beam solution corrected by contributions due to the limited bunch length. In fact the collision of short rectangular bumps with a bunch length of the order of the luminous region length produces less luminosity than the luminosity calculated from the coasting beam approximation.

The luminosity degradation due to non-ideal rectangular line density can be taken into account as a form factor. For bunches whose local longitudinal density \( \lambda(z) \) varies insignificantly along the luminous region, the form factor becomes
\[
\frac{\mathcal{L}}{\mathcal{L}_0} = \int \lambda^2(z) \, dz ,
\]
where the colliding bunches are assumed to be longitudinally equal and symmetric \( \lambda(\Delta z) = \lambda(-\Delta z) \). The crossing of perfect rectangular bunches results in a form factor of unity.

**STRONG BEAM-BEAM LIMITATION FOR RECTANGULAR BUNCHES**

The maximum incoherent tune shift of a single particle in a beam under the influence of the electromagnetic force of a second beam is expressed in both transverse planes by [7]
\[
\Delta Q_x = -\frac{1}{4\pi} \int_{-l_{\text{det}}/2}^{l_{\text{det}}/2} \frac{1}{2m_0 \gamma^2 c^2} \frac{dF_x}{dy} \bigg|_{x=y=0} \beta_x(z) \, dz ,
\]
\[
\Delta Q_y = -\frac{1}{4\pi} \int_{-l_{\text{det}}/2}^{l_{\text{det}}/2} \frac{1}{2m_0 \gamma^2 c^2} \frac{dF_y}{dx} \bigg|_{x=y=0} \beta_y(z) \, dz ,
\]
where \( l \) is the length along which both beams interact without shielding in-between.

The Lorentz force of a single beam in the coordinate system co-aligned with the beam itself can be written as
\[
|F_L| = F_r(r) = \frac{\lambda e^2}{2\pi \epsilon_0 \epsilon r} (1 + \beta^2) \left[ 1 - e^{-r^2/(2\sigma^2)} \right] .
\]

Transforming to the coordinate system co-aligned with the second, counter-rotating beam, which is tilted by the crossing angle \( \theta \) according to the rotation introduced in [7], the beam-beam force from a round, rectangular bunch becomes
\[
F_x(x, y, z) = 2m_0 \beta^2 e^2 \lambda r_p (1 + \cos \theta) \times \zeta(z + \cos \theta + x \sin \theta) \frac{x \cos \theta - z \sin \theta}{(x \cos \theta - z \sin \theta)^2 + y^2} \times \left( 1 - \exp \left[ -\frac{(x \cos \theta - z \sin \theta)^2 + y^2}{2\sigma^2} \right] \right)
\]
and
\[
F_y(x, y, z) = 2m_0 \beta^2 e^2 \lambda r_p (1 + \cos \theta) \times \zeta(z + \cos \theta + x \sin \theta) \frac{y}{(x \cos \theta - z \sin \theta)^2 + y^2} \times \left( 1 - \exp \left[ -\frac{(x \cos \theta - z \sin \theta)^2 + y^2}{2\sigma^2} \right] \right) ,
\]
where \( \lambda = N/l_b \) denotes the longitudinal line density and \( r_p = 1/(4\pi \epsilon_0 \epsilon^2) e^2/m_p \approx 1.535 \cdot 10^{-18} \text{ m} \) is the classical proton radius.
Inserting the Lorentz force into Eqs. (9) and (10) gives the horizontal

\[
\Delta Q_x = \frac{\lambda r_p \beta^*}{2\pi \gamma} (1 + \cos \theta) \left[ \cos \theta \int_{-l/2}^{l/2} \left( 1 + \frac{z^2}{\beta^*} \right) \right. \\
\left. \times \zeta(z + z \cos \theta) \left\{ \frac{1}{z^2 \sin^2 \theta} \left[ 1 - \exp \left( \frac{-z^2 \sin^2 \theta}{2\sigma(z)^2} \right) \right] \\
- \frac{1}{\sigma(z)^2} \exp \left( \frac{-z^2 \sin^2 \theta}{2\sigma(z)^2} \right) \right\} dz \right]
\]

and vertical

\[
\Delta Q_y = \frac{\lambda r_p \beta^*}{2\pi \gamma} (1 + \cos \theta) \left[ \cos \theta \int_{-l/2}^{l/2} \left( 1 + \frac{z^2}{\beta^*} \right) \right. \\
\left. \times \zeta(z + z \cos \theta) \frac{1}{z^2 \sin^2 \theta} \left[ 1 - \exp \left( \frac{-z^2 \sin^2 \theta}{2\sigma(z)^2} \right) \right] dz \right]
\]

beam-beam tune shift. The second term in the horizontal tune shift contributes as a small edge term only if the ends of the bunch are within the integration range and can be neglected. For long bunches, with \( \zeta(z + z \cos \theta) = 1 \) all along the interaction region, this result reduces to the well known tune shift of a coasting beam [5].

The total beam-beam tune spread caused by two alternate horizontal/vertical beam crossings (Fig. 2) is just given by the sum of \( \Delta Q_x \) and \( \Delta Q_y \) [7]:

\[
\Delta Q_{tot} = \frac{\lambda r_p \beta^*}{2\pi \gamma} (1 + \cos \theta) \left[ \cos \theta \int_{-l/2}^{l/2} \left( 1 + \frac{z^2}{\beta^*} \right) \right. \\
\left. \times \zeta(z + z \cos \theta) \frac{1}{z^2 \sin^2 \theta} \left[ 1 - \exp \left( \frac{-z^2 \sin^2 \theta}{2\sigma(z)^2} \right) \right] \\
- \frac{1}{\sigma(z)^2} \exp \left( \frac{-z^2 \sin^2 \theta}{2\sigma(z)^2} \right) \zeta(z + z \cos \theta) dz .
\]

**OPTIMUM LUMINOSITY OF LONG BUNCHES**

In the limit of bunches much longer than the luminous region, the expressions derived above can be approximated so that luminosity and total beam-beam tune spread reduce to [4]

\[
L = \frac{f_0 \lambda h \lambda^2}{\sqrt{\pi}} \frac{1}{\sigma^* \theta}
\]

and

\[
\Delta Q_{tot} = -\sqrt{\frac{2}{\pi}} \frac{\lambda r_p \beta^*}{\gamma \sigma^* \theta}.
\]

Clearly, the luminosity is directly proportional to the bunch length as the average current also increases proportionally when a larger fraction of the ring is occupied. It is interesting to note that both parameters have the same dependence on \( \theta \) so that the luminosity in terms of beam-beam tune shift remains constant for arbitrary crossing angles.

To obtain the conditions for maximum luminosity, two restrictions are assumed: Firstly, the average intensity in the LHC is limited by cryogenic cooling capacities with respect to synchrotron radiation losses or the power capability of the beam dump. Secondly, its peak intensity is ultimately restricted by the strong beam-beam tune spread. As shown in Fig. 3, operation at optimum parameters is achieved when both limitations are reached simultaneously (marked with a black dot). Form Eqs. (17) and (18) the maximum luminosity becomes

\[
L_{max} = \frac{1}{\sqrt{2}} \frac{f_0 \gamma}{r_p \beta^*} N_{max} \Delta Q_{tot},
\]
so that the total beam length to achieve this luminosity is written as

$$n_b l_b = \sqrt{\frac{2}{\pi}} \frac{r_p \beta \gamma^2}{\sigma_\theta N_{\text{max}}} \frac{\Delta Q_{\text{tot}}}{\Delta Q_x \Delta Q_y}.$$  

(20)

The number of equal bunches is defined by \(n_b\).

This result shows that if the approximations to apply Eqs. (17) and (18) are fulfilled and the fraction of particles in the tails can be neglected, the maximum luminosity depends only on the total beam length \(n_b l_b\).

**FLAT AND SUPER-BUNCH UPGRADE OPTIONS**

Based on the luminosity formulae derived above, three flat or super-bunch upgrade schemes are analyzed: Firstly, colliding short and flat bunches by adding higher harmonics to the fundamental RF frequency emulating a barrier bucket like RF amplitude. Secondly, either 16 or 32 nearly nominal bunches are combined to a high intensity long and flat bunch held by multiples of 40.08 MHz. The resulting bunch length is in the order of a few meters. Thirdly, all particles are confined to one single super-bunch per beam held by barrier buckets.

**Short Rectangular Bunches**

Adding multiples of 400.8 MHz to the RF amplitude generated by the superconducting accelerating cavities allows to approach rectangular bunches. The longitudinal phase space including the line density of a bunch held by two and three RF systems is illustrated in Fig. 4 and 5. Even for only three harmonics at \(n \cdot 400.8, n = 1 \ldots 3\) more than 80% of the particles are within the region of homogeneous line density.

It is important to to point out that the simplified formulae Eqs. (17) and (18) cannot be used to calculate the performance of such a scheme in the LHC for crossing angles of some 1 mrad as the bunch length (\(\approx 20\) cm) is of the same order of magnitude as the strong beam-beam interaction length. The beam-beam behaviour is similar to that of the nominal Gaussian bunches and no compensation of the tune shift in the crossing plane occurs so that \(\Delta Q_x \simeq \Delta Q_y\).

The maximum achievable luminosity at a constant beam-beam tune shift of \(\Delta Q_{\text{tot}} = -0.01\) for two alternate interactions points for short rectangular as well as Gaussian bunches versus bunch length is plotted in Fig. 6. Clearly, though short rectangular bunches can produce slightly more luminosity for a fixed intensity per bunch, their potential luminosity at the beam-beam limit may even remain below that of a Gaussian bunch crossing.

Increasing the crossing angle to some 5 mrad would reduce the ratio of bunch to strong interaction length so that rectangular bunches could deliver \(\sqrt{2}\) times more luminosity than Gaussian ones. The beam-beam limit is pushed far away from any realistic beam current, however according to Eq. (17) an average beam current in the range of

\(I_0 > 20\) A is required to profit from luminosity improvements at such large crossing angles. These intensities are not compatible with any present or upgrade schemes of the LHC.

**Long and Flat Bunches**

For bunches of ten times the nominal bunch length a luminosity gain of nearly \(\sqrt{2}\) is nearly reached for a cross-
ing a crossing angle of 1 mrad (see Fig. 7). Combining 16

![Diagram](image)

Figure 7: Comparison of rectangular and Gaussian bunch crossings at various crossing angles (continuous: $\theta = 285 \mu$rad, dashed: 500 $\mu$rad, dot-dashed: 1 mrad, dotted: 5 mrad) in the LHC. The theoretical limitation of the luminosity gain is $\sqrt{2}$ (dashed straight line).

or 32 nearly nominal LHC bunches held by an RF voltage composed of harmonics of $n \cdot 40.08$ MHz could provide such long and flat bunches. The approximation of the ideal rectangular bunch form is equivalent shown above (Figs. 4 and 5).

The maximum obtainable luminosity for the crossing of long and flat bunches is summarized in Tab. 1. The 16 bunch option can reach the double luminosity as the collision of combinations of 32 bunches but at the cost of twice the average beam current.

### Super-Bunches

The third flat bunch upgrade options is discussed in [2]: All particles circulate in a single, sausage like super-bunch per beam which is confined by barrier buckets as illustrated in Fig. 8. According to Eq. (17) the attainable luminosity from the collision of such bunches is proportional to $\propto \lambda^2 l_b / \theta$. In fact, the collision of superbunches is equivalent to the collision of two coasting beams with the difference that only a fraction of the circumference is occupied by particles. As the barrier RF pulses needed to confine the bunches are independent, the bunch length and thus the longitudinal particle density can by easily adjusted to the beam-beam limit by adiabatically moving the phases of the single sinusoidal pulses with respect to each other.

Assuming the same crossing parameters as for the intermediate long and flat bunches, namely $\theta = 1$ mrad a luminosity as presented in Tab. 2 is within reach. It is important to point out that synchrotron radiation must be compensated by an additional dedicated pulsed RF system. The average energy loss of 6.71 keV per turn causes a strong asymmetry of the bunches in the barrier buckets otherwise.

### Summary of the Rectangular Bunch Options

A comparison of the parameters of the three flat bunch options and their impact on the required upgrades of the accelerator and the detectors is given in Tab. 3. Colliding short rectangular bunches at very large crossing angles does not deliver sufficient luminosity with reasonable average beam currents in the LHC. However, the potential performance at the beam-beam limit would be almost two orders of magnitude above the luminosity of the nominal LHC scheme. The super-bunch option represents a flexible scheme as the longitudinal particle density can be adjusted arbitrarily by variation of the barrier phases. This scheme requires two complex barrier bucket RF systems for confinement of the long bunches and the compensation of synchrotron radiation. No realistic scheme to combine all particles to a single long bunch per beam has been found so far but a major breakthrough cannot be excluded.

Comparable performance with respect to luminosity can be achieved by the operation of the LHC with long and flat bunches of intermediate length as summarized in the middle row of Tab. 3. A possible RF manipulation procedure

![Diagram](image)
Table 3: Comparison of the flat bunch luminosity upgrade options in the LHC sorted by ascending bunch length.

<table>
<thead>
<tr>
<th></th>
<th>short and flat bunch</th>
<th>long and flat bunch</th>
<th>superbunch</th>
</tr>
</thead>
<tbody>
<tr>
<td>Total bunch length, $l_b$</td>
<td>40...50 cm</td>
<td>4...5 m</td>
<td>200...400 m</td>
</tr>
<tr>
<td>Total beam intensity for $L_{tot} = 10^{35}$ cm$^{-2}$s$^{-1}$</td>
<td>enormous</td>
<td>intermediate</td>
<td>extensive</td>
</tr>
<tr>
<td>Additional RF systems</td>
<td>800/1200 MHz</td>
<td>40/80/120 MHz</td>
<td>Barrier bucket RF at $\simeq 10$ MHz</td>
</tr>
<tr>
<td>Consequences for physics experiments</td>
<td>limited</td>
<td>significant/difficult modifications</td>
<td>extensive/impossible modifications</td>
</tr>
<tr>
<td>Consequences for the accelerator hardware</td>
<td>limited</td>
<td>significant modifications</td>
<td>extensive modifications</td>
</tr>
<tr>
<td>RF gymnastics to generate long bunches</td>
<td>easy</td>
<td>sophisticated</td>
<td>presently no realistic scheme</td>
</tr>
<tr>
<td>Synchrotron radiation compensation</td>
<td>unnecessary</td>
<td>unnecessary</td>
<td>necessary and difficult</td>
</tr>
</tbody>
</table>

to generate such bunches is therefore presented in what follows.

**GENERATION OF LONG BUNCHES WITH INTERMEDIATE LENGTH**

The generation of long bunches in the LHC, each composed of 16 or 32 nearly nominal bunches, can be achieved by an RF manipulation scheme which starts during ejection from the CERN Super Proton Synchrotron (SPS). The main parts of the scheme are performed at the flat-bottom in the LHC because beam losses at collision energy due to RF gymnastics should be avoided [8]:

1. Bunch lengthening by bunch rotation and matched injection into the LHC.
2. Batch compression and bunch pair merging of $16 \rightarrow 1$ bunches during the flat-bottom in the LHC.
3. Acceleration to 7 TeV collision energy at 40.08 MHz.
4. Final batch compression and formation of the long and flat bunches at the flat-top.

An overview of the scheme of combining 16 bunches is given in Fig. 10. For the combination of 32 bunches an additional sequence of batch compression and bunch merging is prepended. The total duration of the complete procedure is about one minute. For the combination of 16 bunches two tunable RF systems in the frequency range of 40–80 MHz with delivering a voltage of 1.5 MV each are required. In the case of 32 bunch batches, twice the longitudinal emittance must be handled corresponding to 6 MV.

**Beam Transfer from SPS to LHC**

The 400.8 MHz superconducting cavities would generate intolerably large voltages due to transient beam loading so that they have to be removed for the operation with long and flat bunches and the beam coming from the SPS (harmonic number, $h = 4620$) is injected directly into 40.08 MHz ($h = 3564$) buckets in the LHC. The condition for a longitudinally matched transfer (index 1: SPS, index 2: LHC) can be written as

$$\frac{V_1}{V_2} = \left( \frac{R_1}{R_2} \right)^2 \frac{\eta_2 h_2}{\eta_1 h_1} \simeq 0.38 \frac{h_2}{h_1},$$

(21)

where $V$ denotes the RF voltage and $\eta$ is the phase slip factor. The RF voltage $V_1$ can vary from 3 to 8 MV.

If only 3 MV are available at 40.08 MHz as foreseen for the 16 bunch combination scheme, the corresponding RF voltage from Eq. (21) in the SPS would lead to insufficient bucket area to keep the 0.6–0.8 eV bunches within the linear region of the buckets. Therefore, the bunches have to be stretched at the unstable fixed point [9] by switching the RF phase (8 MV/200 MHz) for some 20 turns. Changing the RF phase back to the stable fixed point excites a quadrupole oscillation as the stretched bunch is not matched to the bucket. At the instant of largest bunch length ($\simeq 2.6$ ns) the beam is ejected within one turn and transferred to the LHC.

To perform the combination of 32 bunches a total RF voltage 12 MV in the LHC is demanded anyway and a matched transfer should be possible without previous lengthening in the SPS. However, beam longitudinal stability in the LHC may limit the maximum applicable RF voltage during injection so that the scheme described above is also required.

**RF Gymnastics at Flat-Bottom**

When the injection is completed, each LHC ring will be filled (excluding several kicker and beam dump gaps) with short batches of 16 bunches with gaps of 2 empty bucket
positions \((16 \otimes b \oplus 2 \otimes e)\) respectively 32 bunches with gaps of 4 empty bucket positions \((32 \otimes b \oplus 4 \otimes e)\) in-between. The bunch spacing of 25 ns \((h = 3564)\) corresponds to the nominal spacing in the LHC.

Batch compression [10] and bunch pair merging [11] are the two main ingredients of the subsequent RF manipulation to combine each sub-batch to a single long and flat high intensity bunch. Due to insufficient bucket area, the scheme cannot be completed at injection energy, but the bunches are combined to two dense bunches each containing 8 or 16 initial bunches.

Batch compression is based on the fact that bunches kept by an RF system at \(h_1\) can be transferred adiabatically to a second RF system at slightly higher harmonic \(h_2\), typically \(h_2 = h_1 + 1\) or \(h_2 = h_1 + 2\) if \(h_1, h_2\) denote the harmonic number with respect to the batch. This harmonic hand-over can be performed by increasing the amplitude at the higher harmonic while decreasing the lower harmonic amplitude simultaneously. The compression factor of the batch corresponds to the ratio of the RF wavelengths: \(h_2/h_1\). Literally, \(h_2 - h_1\) buckets are inserted by such a procedure. The adiabatic hand-overs between RF systems are applied repetitively so that the RF harmonic is increased from \(h = 3564\) to 7128. The increment depends on the batch length, \(16\) + 2 or \(32\) + 4 bunch positions, and is chosen as \(\Delta h = 3564/18 = 198\) (2.2 MHz) respectively \(3564/36 = 99\) (1.1 MHz). After the batch compression the bunch spacing is halved while leaving the number of bunches unchanged. The complementary procedure, bunch pair merging, doubles the bunch spacing while reducing the number of bunches by factor. The higher harmonic amplitude at \(h = 7128\) is adiabatically reduced and the RF voltage at \(h = 3564\) is slowly turned on. Applying such a sequence (see Figs. 9 and 10 for the combination of 16 bunches) of batch compression and bunch pair merging fi-

\[
\begin{align*}
\begin{array}{lll}
\text{n}_{\text{batch}} & \text{Pattern} & h & \text{Merge} \\
18 & 16 \otimes b \oplus 2 \otimes e & 3564 & \overset{+198}{\text{3564}} \\
19 & 16 \otimes b \oplus 3 \otimes e & 3762 & \overset{+198}{\text{}} \\
20 & 16 \otimes b \oplus 4 \otimes e & 3960 & /\!\!/2 \\
\vdots & \vdots & \vdots & \vdots \\
36 & 16 \otimes b \oplus 20 \otimes e & 7128 & 7128 \\
\end{array}
\end{align*}
\]

Figure 9: Schedule of harmonic numbers and bunch patterns during long bunch combination of 16 initial bunches.

nally halves the number of bunches. Repetitive application allows to combine the batch sequentially and the bunch pattern develops as \(16 \otimes b \oplus 2 \otimes e \rightarrow 8 \otimes b \oplus 10 \otimes e \rightarrow 4 \otimes b \oplus 14 \otimes e \rightarrow 2 \otimes b \oplus 16 \otimes e\) (16 bunches) or \(32 \otimes b \oplus 4 \otimes e \rightarrow 16 \otimes b \oplus 20 \otimes e \rightarrow 8 \otimes b \oplus 28 \otimes e \rightarrow 4 \otimes b \oplus 32 \otimes e \rightarrow \cdots \rightarrow 2 \otimes b \oplus 34 \otimes e\) (32 bunches).

\[
\Delta E \propto \Delta E(\phi)^2
\]

Figure 10: Mountain range plot of the separatrices (colour scale \(\Delta E(\phi)^2\)) during formation of the long bunches.

4 \(\otimes b \oplus 32 \otimes e \rightarrow 2 \otimes b \oplus 34 \otimes e\) (32 bunches).

**Acceleration**

Due to the small ramp rate of the superconducting magnets in the LHC the acceleration takes place within 20 min at a moderate average energy gain of 0.485 MV per turn. The minimal bucket area of 26.2 eVs for 3 MV (twice that area for 12 MV) RF voltage at 40.08 MHz occurs about one minute after the start of the acceleration and offers a comfortable margin with respect to the expected longitudinal emittance of the bunches (see Tab. 4 for the 16 bunch scheme).

**Final Formation of Long Bunches**

Due to insufficient bucket area at low energy, the final batch compression of the two bunch batches to \(h = 7128\) must take place after acceleration at collision energy. The final formation of the long and flat bunches resembles a bunch pair merging halted in the middle of the process. Starting from the single harmonic RF system, the amplitudes at multiples of 40.08 MHz are adiabatically changed to their final valued needed to hold a bunch with optimized flatness of the homogeneous region as shown in Fig. 4 and 5.
Table 4: Emittance development during the combination of 16 bunches to a flat bunch (16 $\otimes b$ $\otimes 2 \otimes e$).

<table>
<thead>
<tr>
<th>RF parameters</th>
<th>Energy, $E$</th>
<th>Emittance, $\varepsilon_L$</th>
<th>Blow-up</th>
</tr>
</thead>
<tbody>
<tr>
<td>SPS ejection</td>
<td>450 GeV</td>
<td>8 MV at 200 MHz (SPS)</td>
<td>0.8 eVs</td>
</tr>
<tr>
<td>Bunch rotation and LHC injection</td>
<td>3 MV at 40 MHz (LHC)</td>
<td>1.0 eVs</td>
<td>10%</td>
</tr>
<tr>
<td>Blow-up by 400 MHz RF system</td>
<td>3 MV at 40 MHz</td>
<td>1.1 eVs</td>
<td>5%</td>
</tr>
<tr>
<td>Batch compression to two bunches</td>
<td>2 $\times$ 1.5 MV at 40...80 MHz</td>
<td>12.3 eVs</td>
<td>40%</td>
</tr>
<tr>
<td>Acceleration to flat-top</td>
<td>3 MV at 40 MHz</td>
<td>12.9 eVs</td>
<td>5%</td>
</tr>
<tr>
<td>Final formation of the long bunch</td>
<td>2 $\times$ 1.5 MV at 40...80 MHz</td>
<td>28.5 eVs</td>
<td>10%</td>
</tr>
<tr>
<td>Collision mode with long bunches</td>
<td>0.8/0.8/0.4 MV at 40/80/120 MHz</td>
<td>28.5 eVs</td>
<td></td>
</tr>
</tbody>
</table>

LONGITUDINAL BEAM STABILITY

According to the analysis in [12], the beam is most prone to coupled bunch instabilities during the injection procedure (bunches are much smaller than buckets, meaning small synchrotron frequency spread, $\varepsilon_L/\varepsilon_{bucket} \propto \Delta \omega_S/\omega_S$). The synchrotron frequency spread is at least one order of magnitude smaller than for a nominal bunch leading to loss of Landau damping due to the estimated broad band impedance [13] and coupled bunch mode excitation by any kind of narrow band impedance. A higher harmonic RF voltage around $80\,\text{MHz}$ to increase $\Delta \omega_S/\omega_S$ and/or a longitudinal bunch-by-bunch feedback could be used to stabilize the bunches.

It is worth noting that, with respect to the present design of the beam crossing regions [14], the distance between bunches (0.45 $\mu$s/0.90 $\mu$s) is large enough that there are no parasitic beam-beam interactions between different bunches anymore.

HARDWARE REQUIREMENTS

As the initial bunch rotation in the SPS can be performed with the available 200 MHz RF cavities, the only major hardware upgrade for the intermediate long bunch scenario is in the LHC. Two RF systems providing 1.5 MV each in the frequency range of 40...80 MHz are required. The installation of fixed frequency RF systems at 120/160 MHz for an improved line density distribution could be considered and further harmonics may be added later if necessary. Assuming that some 300 kV can be obtained per cavity, the two frequency variable RF systems would consist of 5 (16 bunch scheme) respectively 20 (32 bunch scheme) cavities each. The cavity tuning is foreseen to be done with mechanical elements near the gap, similar to the pneumatic gap short circuits already operational in the PS at CERN [15, 16].

CONCLUSIONS

Analytical formulae for the calculation of luminosity and strong beam-beam limitation of rectangular bunches of arbitrary length are derived and applied to different flat bunch upgrade options for the LHC. It turns out that crossing flat bunches of intermediate length can deliver almost as much luminosity as the crossing of super-bunches avoiding several of their inherent disadvantages.

An RF manipulation scheme to combine batches of 16 or 32 initial bunches to these flat bunches is presented. Though sophisticated, it relies heavily on conventional RF gymnastics which are well proven to work up to highest beam intensities in smaller accelerators and the additional RF hardware is of reasonable size and performance.

ACKNOWLEDGEMENTS

We would like to thank Elena Shaposhnikova, Steven Hancock, Joachim Tückmantel and Frank Zimmermann for the stimulating discussions. One of the authors (H.D.) is grateful to GSI, Darmstadt for supporting the studies at CERN.

REFERENCES


Machine-detector interface and event pile-up: super-bunches versus normal bunches

S. Tapprogge*, Institut f. Physik, Johannes-Gutenberg Univ. Mainz, D 55099 Mainz, Germany

Abstract
A future upgrade of the LHC to achieve peak luminosities larger by up to an order of magnitude compared to the present design value will allow to enlarge the physics potential sizeably. It will have to be accompanied by corresponding upgrades of the ATLAS and CMS detectors, in order to be able to fully exploit this increase in physics potential.

The suitability of different LHC upgrade scenarios is investigated from the experiments (ATLAS and CMS) point-of-view in this contribution. An essential figure-of-merit for these considerations is the number of quasi-simultaneous inelastic proton-proton interactions, occurring during the collisions of two proton bunches.

INTRODUCTION
An upgrade of the LHC machine to achieve a peak luminosity larger by up to one order of magnitude than the design value would allow to extend the LHC physics potential. In order to profit fully from such a machine upgrade, also the ATLAS and CMS detectors would have to be upgraded, retaining as much as possible the capabilities of the present design of the detectors. It has to be kept in mind that the present design is already a challenge and that a further increase in luminosity will be even more challenging to implement from the detector point-of-view. At design luminosity, each bunch collision (occurring with a frequency of 40 MHz) will contain on average 20 — 25 proton-proton interactions. This implies that in each bunch crossing there will be about 750 charged particles (with an average transverse momentum of 0.5 GeV) inside the acceptance (|$\eta|$ < 2.5) of the tracking systems of ATLAS and CMS. For the even higher luminosities of an upgrade, the particle multiplicities will increase even further, as described below.

Details of the upgrade of the detector components depend on the scenario of the machine upgrade, where a range of options is being discussed. Crucial parameters of the upgrade scenarios for the consideration of experimental upgrades are the length of the bunches and the spacing between these, as discussed in this contribution.

Firstly, a brief overview of the physics motivation for an upgrade and the expected increase in the physics potential is given. Next, a brief summary of the upgrade scenarios for the machine as considered in this study is given. Then more details on issues relevant for the various components (tracking, calorimeter, muon system, trigger and data acquisition) of ATLAS and CMS are presented, before the conclusions are drawn.

PHYSICS MOTIVATION FOR UPGRADE
An increase of the peak luminosity to values of $10^{36}$ cm$^{-2}$s$^{-1}$ would allow to extend the physics reach of the LHC in two ways: firstly, the discovery reach would be increased to cover the production of heavier objects and to give access to rare decays, not being within the statistical reach of the nominal LHC luminosity. At the same time, an increased luminosity would allow to improve several precision measurements (of Standard Model processes and also of new physics processes), profiting from the increase in statistics. The following two sections briefly describe some of the key areas where improvements are expected, and are followed by a list of requirements on the performance of an upgraded detector. More details on the physics reach can be found in [1, 2]. Typically, an integrated luminosity of 3000 fb$^{-1}$ per experiment has been assumed in these studies.

Extension of discovery reach
An increase in the peak luminosity by one order of magnitude (assuming the corresponding increase in the integrated luminosity as well) would allow to extend the discovery reach for various scenarios, some of which are listed below:

- heavy Higgs boson (e.g. in the MSSM),
- supersymmetric particles,
- new heavy gauge bosons,
- extra-dimensions,
- quark substructure.

The upper limit on the accessible mass of new objects would increase in general by about 20 — 30%. As a concrete example, in Fig. 1 the discovery reach for squarks and gluinos is shown. The reach is obtained within the mSUGRA model, using final state signatures of jets and large missing transverse energy. The reach in the mass of squarks and gluinos of 2.5 TeV for the LHC could be extended to about 3 TeV for an upgraded LHC.

Also in the case of the search for new heavy gauge bosons the reach in mass increases at an upgraded LHC, as shown in Fig. 2. Assuming that a discovery would be made with at least 10 events, the discovery reach for a Z$'$ boson increases from a mass of 5.3 TeV for the nominal LHC to 6.5 TeV at an upgraded LHC.
Improvements in precision measurements

A sizeable increase in statistics would allow to improve precision measurements of many parameters, as those are often limited in statistics even for the nominal luminosity of LHC. The list of processes where substantial gain in knowledge could be obtained includes the following items:

- couplings of the Higgs boson to fermions and gauge bosons,
- possibly a measurement of the Higgs self-coupling,
- rare decays of the top quark via flavour changing neutral currents,
- triple and quartic gauge boson couplings,
- mass measurements of SUSY particles.

A first concrete example of the improvement in precision measurements is the determination of the Higgs boson couplings, as shown in Fig. 3. An upgraded LHC should allow the determination of the ratios of the couplings to fermions and bosons to about 10% in most cases, which is an improvement by a factor of 2 over the nominal LHC performance.

Requirements on detector performance

In order to profit from the projected increase in statistics due to an upgrade, it is essential that upgraded detectors will have a similar performance as foreseen for the present design. Although some discoveries (beyond the nominal LHC reach) might be possible by using information from calorimetry and muon system alone, it is clear that a full physics coverage and the understanding of new phenomena will require also good tracking capabilities. This is necessary to identify and reconstruct objects such as high-\(p_T\) electrons, photons, hadronic tau decays and to tag the flavour of jets (\(b\)-tagging). All of these signatures rely on good tracking capabilities both for the identification of signal signatures and the rejection of fake signatures from background processes.

UPGRADE SCENARIOS CONSIDERED

Details on the possible implementation of various upgrade scenarios can be found in [3]. In the following, a brief summary of the key scenarios which have been con-
considered in this study is given. For an upgrade of the LHC, three phases are distinguished. In all of these phases and scenarios, it is assumed that the beams collide only in two interactions regions.

**Phase 0:** aims at stretching the machine performance to its limits. This could be achieved essentially without changes to the LHC hardware, but might imply an upgrade of the injectors. The peak luminosity would increase by a factor of up to 2.3 compared to the design value of $10^{33} \text{ cm}^{-2}\text{s}^{-1}$.

**Phase 1:** would include the upgrade of the interactions regions, to achieve stronger focusing of the beams at the interaction points. Various scenarios have been developed and the resulting boundary conditions for the upgrade of the experiments are described below.

**Phase 2:** would aim at an increase in the center-of-mass energy for the $pp$ collisions as well. An upgrade to a value of $\sqrt{s} = 28$ TeV would imply the development of new dipoles with a bending field of 15 T, possibly as well an upgrade of the SPS (with new superconducting magnets) to inject protons at 1 TeV.

Only phase 1 has been considered in detail in this document. For a fixed value of the peak luminosity, the more demanding and expensive upgrade in center-of-mass energy could easily be handled by the detectors in their nominal design, as the number of particles (and the number of simultaneous inelastic events per bunch crossing) only increases slowly (logarithmically) with center-of-mass energy. In case of larger $\sqrt{s}$ values and a luminosity higher by one order of magnitude, the same arguments would apply for phase 2 as these do for phase 1. In the discussion, the stretching the luminosity by a factor of 2 has been included as well. It has to be always kept in mind, that the scenarios listed below should only be taken as guidelines, as details will keep changing.

**Parameters for phase 1 upgrade scenarios**

In Table 1, an overview of the most important parameters for the various scenarios investigated in details is given. The major parameters considered are the peak luminosity value, the bunch spacing and the bunch length, and the expected number of inelastic interactions per bunch crossing. The first scenario in the table concerns the nominal LHC parameters, the second one those expected to be obtainable in phase 0 (ultimate performance) without major hardware changes. The next two scenarios (Piwinski-1 and -2) assume larger crossing angles and bunch lengths (in order to avoid limitations from beam-beam interactions). Also the latter should not involve major hardware changes. The next three scenarios listed however are real phase 1 scenarios and would involve major changes to the interaction regions. The scenario ‘IR-upgrade’ would e.g. require new quadrupoles to achieve stronger focusing at the interaction
Table 1: Overview of the parameters for various machine upgrade scenarios used in this study. $L$ is the peak luminosity (in units of $10^{34} \text{ cm}^{-2}\text{s}^{-1}$), $\Delta T$ the bunch crossing time, $l$ the bunch length and $N_t$ the number of interactions per bunch crossing.

<table>
<thead>
<tr>
<th>Scenario</th>
<th>$L$</th>
<th>$\Delta T$ (ns)</th>
<th>$l$ (cm)</th>
<th>$N_t$</th>
</tr>
</thead>
<tbody>
<tr>
<td>nominal</td>
<td>1.0</td>
<td>25</td>
<td>7.55</td>
<td>23</td>
</tr>
<tr>
<td>ultimate</td>
<td>2.3</td>
<td>25</td>
<td>7.55</td>
<td>53</td>
</tr>
<tr>
<td>Piwinski-1</td>
<td>3.6</td>
<td>25</td>
<td>15.20</td>
<td>83</td>
</tr>
<tr>
<td>Piwinski-2</td>
<td>3.2</td>
<td>75</td>
<td>40.00</td>
<td>221</td>
</tr>
<tr>
<td>IR-upgrade</td>
<td>4.6</td>
<td>25</td>
<td>3.78</td>
<td>108</td>
</tr>
<tr>
<td>Piw-IR-upg.</td>
<td>6.3</td>
<td>75</td>
<td>20.00</td>
<td>435</td>
</tr>
<tr>
<td>super-bunch</td>
<td>9.0</td>
<td>$\approx 88000$</td>
<td>7500.00</td>
<td>$\approx 900000$</td>
</tr>
</tbody>
</table>

point. This could be combined again with the introduction of a larger crossing angle and larger bunch length (Piw-IR-upg.) to achieve a further increase in peak luminosity. The final scenario (super-bunch) would involve the filling of the machine with two very long bunches ($\sigma_z = 75 \text{ m}$) circulating in opposite direction. Those bunches would provide quasi-continuous pp interactions during a period of about 1 $\mu$s (out of the 88 $\mu$s revolution time). Although the bunch length varies significantly between these scenarios, the resulting longitudinal size of the luminous region is quite similar and has a value of $\sigma_z = 2 \sim 4 \text{ cm}[7]$.

As can be clearly seen from the table, the expected number of interactions per bunch crossing strongly varies with the various scenarios. In the following it will be shown, that this parameter is one of the most important figure-of-merits for assessing the impact of a machine upgrade on the detector performance and determining strategies for the upgrade of the detectors. Longer bunches (or longer bunch spacings) are preferred from the machine point-of-view for an upgrade, as these would minimize the impact of the so-called electron cloud effect on the machine performance, the magnitude of which is presently unknown and will have to be determined during the first operation of the LHC.

Boundary conditions for the various scenarios

Based on the given peak luminosity value for each scenario and the bunch length and spacing, an important parameter for the experimental consideration can be derived, namely the number of (quasi-)simultaneous inelastic interactions per bunch crossing (or during a fixed time interval for the super-bunch scenario).

Specifically for the super-bunch scenario, where for a duration of about 1 $\mu$s almost continuous collisions will be provided (out of the 88 $\mu$s revolution period of a bunch) in collisions of bunches with a length of 300 $\mu$m, the following conditions have to be considered. The expected interaction rate for a luminosity of $10^{35} \text{ cm}^{-2}\text{s}^{-1}$ is about 10 GHz. Thus for the collision of two such super-bunches, $10^6$ interactions will occur during 1 $\mu$s. This gives a rate of 1000 interactions per 1 ns time interval. During a time interval of 25 ns thus the huge number of about 25000 interactions is obtained, which could be compared to about 25 for the present bunch spacing of 25 ns at design luminosity or to 125 for a luminosity of $10^{35} \text{ cm}^{-2}\text{s}^{-1}$ and a bunch spacing of 12.5 ns (where it has been assumed that new electronics would be built to match the reduced bunch spacing in an optimal way).

ISSUES FOR THE VARIOUS SUB-SYSTEMS

In the following, the most important issues are described for the major components of the general purpose detectors, ATLAS and CMS. These are the tracking detectors, the calorimeters, the muon systems and the trigger/DAQ systems. For a detailed description of the nominal design of the ATLAS and CMS detectors, the technical proposals provide more information (ATLAS[4], CMS[5]). For ATLAS, also the following reference exists[6].

Tracking detectors

The tracking detectors are the parts most affected by the increase in luminosity, which relates into an increase in the flux of particles by one order of magnitude. For an unchanged geometry and sensor layout and for the same bunch spacing value of 25 ns, this would lead to an increase of the average occupancy by a factor of 10. This would clearly imply losses in efficiency for the pattern recognition and an increase in the rate of fake tracks, which would in turn deteriorate the rejection of fake objects. It is also obvious that also from the radiation hardness point-of-view the detectors close to the beam pipe can not be operated efficiently under these conditions (and an exchange of sensors and electronics mounted on these every year or more often is obviously not feasible). Depending on the value of the bunch spacing and on the length of a bunch, this situation could be improved (e.g. by shorter bunch spacing) or further degraded (very long bunch as quasi-continuous beam).

One has to keep in mind that due to the dimensions of the tracking detectors, typical outer radii of 1 m or more, particles from the next event will enter the tracking detectors before all particles from the previous one have left for very short bunch spacing, as a high energy particle traverses 0.3 m in about 1 ns.

A further issue concerns the desire to minimize inefficiencies, which e.g. in the case of ATLAS should be less than 1%. In order to achieve this in the most demanding system, the pixel detector – being closest to the beam pipe – a maximum value for the deadtime for one pixel has been derived (before this pixel can accept the next hit). The value chosen is 2.5 $\mu$s for the pixel sensors, except for the inner most layer (where it is only 0.5 $\mu$s). Higher particle fluxes would – without changes to the front-end and read-out electronics – clearly lead to larger inefficiencies and further deterioration in the pattern recognition.
Calorimetry

The most important impact for the calorimetry will come from the increase in the number of inelastic interactions occurring for one bunch crossing, which leads to a large contribution from the so-called pile-up noise to the energy resolution. Furthermore, the energy flow originating from these events will change lower limits in transverse energy for reconstructing e.g. jets and will also impact on isolation criteria for the reconstruction of objects such as electrons and photons.

For the case of the ATLAS liquid Argon calorimeters, the ionisation signal is shaped and then sampled in 25 ns intervals. Using a procedure of optimal filtering, the signal amplitude and time are extracted from this set of discrete signal values. For a shorter bunch spacing of e.g. 12.5 ns, the correct identification of the bunch crossing might become difficult, if electronics with the same shaping time as foreseen now and the sampling of 25 ns were to be kept. Shorter shaping times lead however to a larger electronic noise contribution to the energy resolution; whereas the pile-up noise will obviously decrease with shorter shaping times. The shaping time would need to be optimized and it is expected that the overall degradation to the energy resolution will scale with $\sqrt{L}$, i.e. by a factor of about 3 for an increase in luminosity by a factor of 10.

High particle fluxes can lead to the build-up of space-charges in the liquid Argon, which are most important in the forward direction (small angles to the beam pipe where the highest fluxes occur). This build-up will decrease the signal collection efficiency and thus impact the energy measurement. Preliminary studies indicate that for all of the above scenarios (except for the super-bunch one) the critical energy density due to the size of the space charge effect will be reached only for regions with $|y| > 2$ or higher (i.e. in the forward direction). In contrast, in the super-bunch scenario the value expected for the energy density from space charge effects would be above the critical one by one order of magnitude or more over the full acceptance of the calorimeter.

Muon system

The performance of the muon system is a somewhat less critical item. The increase in particle fluxes might necessitate a restriction of the muon acceptance to a region of $|y| < 2$, as an increase in the amount of shielding around the beam pipe is going to be non-trivial. Clearly the increase in occupancies of the muon detectors will impact the efficiency of the pattern recognition, most so (as for the tracking) in case of the super-bunch scenario. In general, the high rate operation and stability will have to be investigated in more detail.

Trigger and data acquisition (TDAQ)

The first level trigger will crucially depend on the value chosen for the bunch spacing in an upgrade. For a bunch spacing of 12.5 ns, it would clearly be preferable to rebuild this system with the capability to operate at 80 MHz. It could be envisaged to keep (parts of) the present system and to identify which of the two bunch crossings actually contains the event leading to the trigger signature. However this would require larger readout bandwidths, in order to cope with a factor of at least two larger data volumes to be transferred from the front-end electronics to readout buffers.

For a superbunch scenario, a completely different approach would have to be taken, as for a period of 1 $\mu$s the system would essentially see a continuous beam. When using a free running clock of e.g. 100 MHz, during the 1 $\mu$s period of collisions a readout rate of 100 kHz would result. Clearly, this scenario would also put much more precise timing requirements on the system.

The other parts of the TDAQ system, as the higher level triggers and the data acquisition, will depend more on the details of the upgraded detector components and the implications on the amount of data produced by the upgraded parts.

General issues

The increase in particle fluxes will imply a sizeable increase in radiation background and activation levels. This needs to be taken into account when designing new detectors, as a replacement each year is not a feasible and affordable option. Also aspects relating to services and integration of new components will need to be taken into account from the very beginning in the planning of a detector upgrade, as these constraints can not easily be changed.

The above issues do not depend very much on the details of the bunch structure of a particular upgrade scenario for the machine.

Changes to the layout of the interaction regions, which might imply a movement of machine elements closer to the interaction point (IP), will have to be carefully considered. This will depend on details of the detector upgrades, e.g. on the continued installation of calorimeters in the forward direction. Moving machine elements closer to the interaction point will also imply that protection absorbers (such as the TAS) will have to be moved closer to the IP. Here one has to worry both about the activation of these elements and possible back-splash from interactions of the collisions products in the absorber. For the case of CMS, a movement of a few m might be feasible, but depends as stated above very much on the forward calorimeter. The presently foreseen access and installation scenarios for ATLAS would most likely not allow movements by distances larger than a few 10 cm, with similar worries about the impact on the forward calorimeter.

CONCLUSIONS

The proposed upgrade of the LHC to peak luminosities of up to $10^{35}$ cm$^{-2}$s$^{-1}$ will need to be complemented by
corresponding upgrades of the multi-purpose detectors, ATLAS and CMS, in order to fully exploit this increase in the physics potential. For such a detector upgrade, it is very important that a similar performance as the nominal one will be achieved. This implies that especially the upgraded detectors must have good tracking capabilities.

Various scenarios have been evaluated in their impact on upgrades of the relevant components of the ATLAS and CMS detectors. Such an upgrade is in any case a non-trivial undertaking and will require a focused and strong R&D effort. A crucial figure-of-merit in the assessment of the impact of various upgrade scenarios is the expected number of proton-proton interactions occurring per bunch crossing (or during a certain time period). Based on the physics motivation of such an upgrade of the LHC luminosity by an order of magnitude, it is not seen how in case of the super-bunch scenario, this increase in luminosity could be exploited by an upgraded ATLAS or CMS detector. Even if a complete rebuild of the detectors would be considered, which is beyond the aim for such a modest upgrade, it is not clear that a viable design could be achieved to perform the relevant measurements.

An important issue to be considered in narrowing down the number of scenarios is the integrated luminosity delivered by a future machine, under stable running conditions. If this were to be reached with a somewhat smaller peak luminosity (and possibly a modest increase in running time by a few 10 %), this would be clearly preferable to the case where the highest peak luminosity values are achieved, however possibly at the expense of poor duty cycles and possibly increased machine related backgrounds.

ACKNOWLEDGEMENTS

The author would like to thank the organizers for arranging this interesting workshop and the invitation to present this contribution. He is also grateful to F. Ruggiero and F. Zimmermann for several clarifying discussions before the workshop. The author would finally like to acknowledge very useful discussions with several of his colleagues in ATLAS and CMS during the preparation of this presentation.

REFERENCES

A MODERN ANSWER IN MATTER OF PRECISION TRACKING:
STEPWISE RAY-TRACING

F. Méot (CEA DAPNIA, Saclay) and F. Lemuet (CERN AB/ABP)

Abstract

Precision tracking is not only a matter of integrator simplicity, it also requires accuracy in modelling of magnetic fields, their non-linearities and possible defects. Stepwise ray-tracing can make the best use of these two crucial prerequisites regarding precision integration: allied with the computing speed of modern computers, this results in high performance tools. The topic is discussed through recent developments in the ray-tracing code Zgoubi, aimed at multturn tracking in the strongly non-linear fields of fixed field alternating gradient synchrotrons.

INTRODUCTION

From the early years of synchrotron developments, stepwise ray-tracing has been considered a good technique to integrate particle motion, allowing the drawing of machine parameters from single- or multi-turn tracking, possibly using magnet field maps [1]. The developments presented here are based on such methods using the ray-tracing code Zgoubi [2].

Zgoubi has long been used in synchrotron studies (cf., LHC [3], FNAL recycler ring [4], proton storage rings [5]), and the recent developments in the code discussed here further permit the difficult simulation of large amplitude, multi-turn 6-D tracking in scaling and non-scaling FFAGs, for which only a few codes have been applied [6].

Ray-tracing thus offers a mean for fast optimization of FFAG magnet geometry and fields as constrained by design parameters; it provides correct simulation of multturn motion, with such outcomes as the right computation of lattice tunes, tune variations, time of flight, etc.; it yields precision 6-D multturn tracking and motion stability limits in FFAGs.

In the following, the methods for simulating FFAG fields are described, and then applied to various problems of 6-D or 4-D tracking in FFAG rings, scaling and isochronous.

RAY-TRACING METHOD, INGREDIENTS

We first recall the ingredients of the Zgoubi method that intervene in the implementation of dipole N-plet simulations.

Position and velocity The integration method is based on stepwise resolution of Lorentz equation by a technique of Taylor series. The working frame is shown in Fig. 1. Position and velocity of a particle subject to \( m \ddot{\mathbf{u}} / dt = q \mathbf{u} \times \mathbf{B} \) are tracked using truncated Taylor expansions in the integration step \( \Delta s \):

\[
\begin{align*}
\bar{R}(M_1) & \approx \bar{R}(M_0) + \bar{u}(M_0) \Delta s + \bar{u}'(M_0) \frac{\Delta s^2}{2} + \ldots \quad (1) \\
\bar{u}(M_1) & \approx \bar{u}(M_0) + \bar{u}'(M_0) \Delta s + \bar{u}''(M_0) \frac{\Delta s^2}{2} + \ldots
\end{align*}
\]

where \( \bar{u} = \hat{u} / v \), \( \Delta s = v \Delta t \), \( \bar{u}' = \frac{d\bar{u}}{d\bar{s}} \), \( m\bar{u}' = q B_0 \bar{u} \), and with the derivatives \( \bar{u}^{(n)} = d^n \bar{u} / d\bar{s}^n \) given by \( \bar{u}' = \bar{u} \times \mathbf{B} \), \( \bar{u}'' = \bar{u}' \times \mathbf{B} + \bar{u} \times \mathbf{B}' \), \( \bar{u}'' = \bar{u}'' \times \mathbf{B} + 2\bar{u}' \times \mathbf{B}' + \bar{u} \times \mathbf{B}'' \), etc.

Taylor coefficients Computation of the coefficients in Eqs. 1 requires the knowledge of the magnetic field \( \mathbf{B}(s) \) and derivatives \( d^n \mathbf{B} / ds^n \) \( (n \leq 5) \) in the orthogonal frame \((O,X,Y,Z)\) (Fig. 1). On the other hand, the magnetic field in a dipole can be obtained from a mid-plane model of the vertical field component (the horizontal component is zero by symmetry), in cylindrical coordinates, of the form \( B_z(r, \theta) = B_{z0} F(r, \theta) \mathcal{R}(r) \), with factors \( F(r, \theta) \) and \( \mathcal{R}(r) \) accounting for the longitudinal (e.g., field fall-offs at dipoles’ ends) and for the transverse (e.g., transverse non-linearities) variation of the dipole field. The way the mid-plane field and its derivatives \( B_z(r, \theta) \), \( \mathcal{R}^{k+l+m} \), at all \((r, \theta)\) are obtained from this model is detailed below.

Once this is done, a transformation from the cylindrical frame of the magnet into the Cartesian frame in Fig. 1 is performed, next, \( Z \)-derivatives and extrapolation off mid-plane are obtained from Maxwell equations and Taylor expansions, thus yielding the 3-D field description

\[
\frac{\partial^{k+l+m} \bar{B}}{\partial X^k \partial Y^l \partial Z^m}
\]

Eventually, \( \bar{B}(s) \) and \( d^n \bar{B} / ds^n \) needed in Eqs. 1 are derived by appropriate coordinate transformations.

STRONGLY NON-LINEAR FIELDS

Two new procedures, named “DIPOLES” and “FFAG”, have been installed in the ray-tracing code for the purposes...
They can account for overlapping fields in the case of neighboring dipoles (Figs. 2-a,b). Dipoles are defined by their parameters as wedge angles, pole curvature, fringe field extents, etc., and are positioned within a sector region with angle $\Delta T$, by means of angles $ACN_i$. The $(r, \theta)$ field dependence has the form

$$B_z(r, \theta) = B_{z0,i} R_i(r) F_i(r, \theta)$$

(2)

(the index $i$ stands for the dipole of concern in an $N$-uplet). “DIPOLES” and “FFAG” differ by the radial behavior, respectively

$$R_i(r) = b_0 + b_1 \frac{r - R_{0,i}}{R_{0,i}} + b_2 \left( \frac{r - R_{0,i}}{R_{0,i}} \right)^2 + \cdots$$

(3)

$$R_i(r) = (r/R_{0,i})^{K_i}$$

(4)

The first form of $R_i(r)$ is proper to simulate FFAG (i.e., $\kappa \geq 0$) with $g_0$ the dipole gap. A dipole having two EFBs (entrance and exit) with each one its own fringe field factor, the resulting form factor at $(r, \theta)$ due to dipole $(i)$ of the $N$-uplet is thus taken to be

$$F_i(r, \theta) = F_{\text{Entrance}}(r, \theta) \times F_{\text{Exit}}(r, \theta)$$

The total mid-plane field and derivatives at $(r, \theta)$ in an $N$-uplet are obtained by summing the contributions of the $N$ dipoles taken separately (e.g., $N = 3$ in Fig. 2), namely

$$B_z(r, \theta) = \sum_{i=1,N} B_{z0,i} R_i(r) F_i(r, \theta)$$

$$\frac{\partial^{k+l}}{\partial r^k \partial \theta^l} \vec{B}_z(r, \theta) = \sum_{i=1,N} \frac{\partial^{k+l} F_i(r, \theta)}{\partial r^k \partial \theta^l}$$

(5)

Eventually, the 6-D field model $\vec{B}(r, \theta, z)$ and derivatives $\frac{\partial^{k+l}}{\partial r^k \partial \theta^l} \vec{B}_z(r, \theta, z)$ are deduced by $z$-extrapolation. Sample $B_z(r, \theta)$ patterns, using the scaling field model in Eq. 4, are given in Figs. 2-b-c, a simulation of the field in an FFAG triplet with characteristics drawn from the KEK 150 MeV proton machine [11].

**6-D TRACKING IN A SCALING FFAG**

We now show that these methods provide correct results, by applying it to 6-D tracking in a scaling FFAG ring.

A 12-cell scaling FFAG ring is considered, representative of the KEK 150 MeV proton FFAG [11]. The cell is a 30 degree sector encompassing a DFD triplet, with $K = 7.6$ in Eq. 4, yielding field on closed orbits as schemed in Figs. 2-b-c, and quasi-zero chromaticity in both planes. Closed orbits in a cell and one-turn tunes are displayed in Fig. 3; with optimized integration step size, tunes values can be guaranteed with good accuracy, better than $10^{-4}$. Other first order results, as drawn from multturn tracking, are displayed in Tab. 1 and show satisfying consistency with published data [11, 12], the momentum compaction satisfies $\alpha \approx 1/(1 + K)$. The vertical chromaticity is not exactly zero due to the fringe fields (zero vertical chromaticity is obtained as expected when a geometrical model with hard edge is used [7]). Fig. 4 shows sample phase space motion at 50 MeV. The horizontal symplecticity is very good. The vertical motion shows confined emittance spread, attributed to non-linear coupling.

Next, stationary bucket dynamics has been simulated (Fig. 5) assuming a single cavity located in a straight section, with peak voltage 19 kV. The agreement with theory
Table 1: First order and longitudinal motion tracking results.

<table>
<thead>
<tr>
<th>E (MeV)</th>
<th>E orbit length (m)</th>
<th>frev (MHz)</th>
<th>mom. synchr. (MeV)</th>
<th>compac. tune</th>
</tr>
</thead>
<tbody>
<tr>
<td>10</td>
<td>28.6333</td>
<td>1.5165</td>
<td>0.11605</td>
<td>0.01133</td>
</tr>
<tr>
<td>22</td>
<td>29.9794</td>
<td>2.1245</td>
<td>0.11611</td>
<td>0.00759</td>
</tr>
<tr>
<td>43</td>
<td>31.1885</td>
<td>2.8089</td>
<td>0.11616</td>
<td>0.00534</td>
</tr>
<tr>
<td>125</td>
<td>33.2724</td>
<td>4.2386</td>
<td>0.11619</td>
<td>0.00291</td>
</tr>
</tbody>
</table>

(e.g., bucket height, synchrotron tunes - Tab. 1) is excellent over the all energy span of the FFAG.

Eventually, a full acceleration cycle, 2 $10^4$ turns from 12 to 150 MeV, using 20 deg. synchronous phase, has been performed, sample results are given in Fig. 5.

**Using field maps**

Magnetic field maps can be used (Fig. 6-a), even in such highly non-linear problem. This is illustrated in Fig. 6-b which displays the horizontal motion stability limits and corresponding large amplitude tunes, at five different energies, as obtained by multturn tracking, which clearly show good simplecticity behavior. There are two *sine qua non* conditions in getting precision, multturn tracking. First, the integrator must be good, RK4 methods for instance would not allow simplectic tracking too far out of the median plane, by contrast with the Zgoubi method. Second, the map mesh must be dense, so as to insure a good interpolation of the - fast oscillating, see Fig. 2-d - derivatives which intervene in Eq. 1.

**COMMENTS**

Computation of field derivatives by numerical differentiation from the mid-plane geometrical field model (Fig. 2) yields good tracking symplecticity, in particular transverse motion (Fig. 4) can be explored up to stability limits. However, using analytical expressions instead for computing the derivatives insures best precision, and allows faster tracking, by a factor of more than 2.

**CPU time**

<table>
<thead>
<tr>
<th>CPU time (seconds per turn per particle)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Pentium III 1 GHz</td>
</tr>
<tr>
<td>2nd order</td>
</tr>
<tr>
<td>4th order</td>
</tr>
</tbody>
</table>

**Figure 4:** (a) : pure radial motion, particles launched with $r_0 = r_{c.o.} + 0.5$ cm (1) and at the stability limit (2). (b) : vertical motion, given $r_0 = r_{c.o.}$.

**Figure 5:** (a) : stationary bucket in the 100 MeV orbit region. (b, c) : respectively $(r, r')$ and $(r, z)$ motions during $12 \rightarrow 150$ MeV acceleration, for a particle launched near the 12 MeV horizontal closed orbit with $z_0 = 1$ cm ; the vertical damping follows $\sqrt{B\rho}$.

**Figure 6:** (a) : magnetic fi eld along closed orbits, drawn by tracking in TOSCA 3-D fi eld maps [13]. (b) : corresponding horizontal stable motion limits and large amplitude tunes.
upon 12 → 150 MeV acceleration in the 12 cell FFAG
ring (conditions as in Fig. 5), using two different proces-
sors, Pentium III 1 GHz or Xeon 2.8 GHz, under Linux
system. An integration step size Δs = 0.5 cm is consid-
ered, derivatives are computed with either the analytical or
the numerical method, up to either second or fourth order
as indicated in Tab. 2. Such computing speed means that
one can envisage overnight runs on computer network sys-
tems, aiming at such goals as long-term DA tracking, 6-D
multi-turn beam transmission, resonance crossing studies.

ISOCHRONOUS LATTICE
A positive chromaticity, isochronous FFAG cell (Fig. 7-
a), has been designed by G. Rees [8], for the purpose of fast
(16-turn) acceleration of muons in the Neutrino Factory.
Isochronism has been obtained by means of the non-linear

transverse field profiles shown in Fig. 7-b,c,d. Such field
shapes can be reproduced for the ray-tracing purposes by
using the mid-plane field model in Eq. 3, given adequate
bₙ coefficients obtained by matching, and yielding typical
fields on closed orbits shown in Fig. 7-e.

A main interest of this particular type of “non-scaling”
FFAG design, with cyclotron-like longitudinal behavior, is
in its yielding best use of the high gradient, 200 MHz RF
acceleration. It also allows use of insertions in the ring.

This type of lattice provides a good illustration of the
power of stepwise ray-tracing : the design parameteriza-
tion requires precision, in particular the isochronism has
to be controlled at a the 10⁻⁸ level, which means neces-
sary accuracy on the description of the strongly non-linear
magnetic fields in the cell dipoles, and on ray-tracing. Typ-
ical outcomes of good symplecticity tracking are the mo-
tion stability limits - in other words, the cell acceptance, of
prime interest - as displayed in Fig. 8.

ADJUSTED FIELD PROFILE LATTICE
This type of non-scaling FFAG lattice has recently been
proposed by A.G. Ruggiero for application in GeV range
proton machines [10, 14].

The longitudinally “Adjusted Field Profile” causes the
index to be a function of the radial displacement x and
of the longitudinal position, that is n = n(x, θ) (Fig. 9),
so cancelling the momentum dependence of the focusing
strength, a consequence being a reduced “non-scaling” :
the variation of the total tune is only of the order of a frac-
tion of an integer over the full acceleration cycle (Fig. 10).
This method has been applied with dipole triplet cells that
have been found to be advantageous, especially in the FDF
configuration that yields low dispersion (Fig. 10). This type
of design is believed to yield competitive technology that
can allow beam performance at the level of the other accel-
or architectures. A main feature is in the compactness of the magnets ensuing from the much reduced beam excursion, compared to scaling FFAG.

Developments in the ray-tracing code Zgoubi are now in progress, in order to allow simulation of these \((x, \theta)\)-dependent, non-linear, sector fields. The principle is in using the polynomial mid-plane modelling of Eq. 2 with radial dependence \(R_i\) as in Eq. 3, yet with the \(b_i\) coefficients functions of the azimuth \(\theta\) in all three dipoles, namely

\[
R_i = 1 - 3(x, \theta) = b_{0,i}(\theta) + b_{1,i}(\theta) x + b_{2,i}(\theta) x^2 + \ldots
\]

Acknowledgements

I thank G. Rees and A. G. Ruggiero for their comments on the manuscript.

REFERENCES


Abstract

For both the SLC damping rings and the DAΦNE collider a systematic approach to understanding single bunch, longitudinal, current dependent behavior was taken: First, using a bunch significantly shorter than nominal, a careful calculation of the wakefield of the entire vacuum chamber was obtained. This “pseudo-Green” function was then used in bunch lengthening and instability calculations. We review, for both projects, the history of these calculations and comparisons with measurement.

INTRODUCTION AND CONCLUSION

In designing an electron storage ring one important consideration is the longitudinal, broad-band impedance of the vacuum chamber, and the effects of this impedance on single bunch behavior of the beam. Possible effects include: potential well bunch shortening, a threshold current, and—above threshold—energy spread increase and bursting (“saw-tooth”) behavior. Any of these behaviors may be deleterious to the performance of a collider, light source, or storage ring, and may need to be avoided.

In particular, for stable, reliable operation it is often desirable that the threshold to the instability be above the operating current. The analysis of stability of a ring normally begins (and often ends) with the Boussard criterion [1]:

$$\frac{\hat{I}[Z(n)/n]}{2\pi\alpha E\sigma_\delta^2} \leq 1,$$

with \(\hat{I}\) the peak current within the bunch, \(Z\) the longitudinal (broad-band) impedance, \(n = \omega_1/\omega_0\) where \(\omega_1\) is a typical bunch frequency and \(\omega_0\) is the revolution frequency, \(\alpha\) is momentum compaction, \(E\) is energy, and \(\sigma_\delta\) is relative energy spread in the beam. Typically, the impedance is calculated for a few important vacuum chamber objects, the contributions are added together, and then inserted into the Boussard equation to estimate the threshold current.

In the studies on the damping rings of the Stanford Linear Collider (SLC) and the DAΦNE collider a new, more systematic approach was used. In both cases, starting with drawings of the vacuum chamber components, an accurate wakefield representing the entire ring was numerically obtained. The driving bunch in the calculations was only a fraction in length of the nominal bunch length, allowing one to use the wake as a pseudo-Green function in subsequent potential well and instability calculations. No adjustable parameters nor fitting was used. In the following two sections we give the history of the calculations and comparisons with measurement for the SLC damping rings and for DAΦNE.

In reading these two sections we see that our approach was quite successful in reproducing measured data (though not perfectly), allowing us to make predictions and obtain insights into the longitudinal, current dependent behavior of the beam. Agreement was found for the current dependence of bunch shape (especially good agreement in the case of DAΦNE), bunch length, synchronous phase; for the threshold current, and—above threshold—the oscillation frequency of the instability. The calculations defined the design strategy of almost all principal vacuum chamber components— in the case of DAΦNE, and led to a redesign of the entire vacuum chamber—in the case of the SLC. Also, through the work of the SLC damping rings a new instability (the weak instability) was discovered, triggering theoretical work to understand it. Note, by the way, that the Boussard criterion often has only an order of magnitude value, and it has nothing to say about the weak instability.

Our approach to impedance calculations appears to have been quite successful; so it is somewhat surprising that it has been seldom used since, and then typically with little success. For example, at the ATF storage ring at KEK the same approach to calculations was followed; nevertheless, bunch length measurements indicate that there is a large amount of still unaccounted for impedance in the ring [2]. It may be that, given the complicated nature of some vacuum chamber components, calculations as described here can still be somewhat of an art form, and one not guaranteed of easy success.

The subject of this report was to be “Impedance Codes and Benchmarking.” By “benchmarking” one can, for example, mean comparing the results of programs, comparing impedance calculations with bench measurements, or (the meaning we pursue) the calculation of ring wakes and consequent current dependent behavior, and then comparison with measurement in the ring. Three types of codes are used in such analysis: impedance (wakefield) calculation codes, threshold finding through a perturbation solution of the Vlasov equation, and tracking for studying behavior above threshold. The programs used (described in the text) seem to have been up to the task for the SLC damping rings and for DAΦNE. However, for all three categories there are recent improvements that should aid the next generation of designers.

In wakefield calculations with short bunches and long structures, the so-called “mesh dispersion” can result in totally wrong results when straightforward calculation is pur-
su ed. A. Novokhatski, et al., for cylindrically symmetric structures [3], and more recently Zagorodnov and Weiland, for 3D structures [4], have developed methods to alleviate this problem. The Oide Vlasov equation solver gives a forest of stable modes (an artifact) amidst any unstable mode one might be searching for [5]. R. Warnock, et al., have reformulated the problem to avoid this artifact [6]. Also, the Oide method, for some impedances, simply fails; the new method seems to have more chance of success (though it can also fail) [7].

As for studying behavior above threshold, R. Warnock and J. Ellison have recently developed a program that solves the Vlasov-Fokker-Planck (VFP) equation, that is more accurate than simple macro-particle tracking [8]. It can, for example, solve also for a coherent synchrotron radiation (CSR) driven, microbunching instability (it appears that such an instability has been observed, e.g. at BESSY II [9]), which is likely impossible to treat with simple tracking. Comparison with measurements of the saw-tooth behavior in the SLC (described in the text) shows promise that one can explore even such complicated behavior numerically with some accuracy.

In the Appendix of this report we present our Panel Comments given at the CARE-HHH workshop held at CERN in November 2004.

SLC DAMPING RINGS

Introduction

In the Stanford Linear Collider (SLC) the beam, after leaving the gun, was stored a few damping times in a damping ring, compressed in the ring to linac (RTL) transfer line, accelerated in the linac, turned around in the arcs, and collided with the opposing beam in the interaction region. Soon after commissioning of the collider it was realized that, in the damping rings, the threshold (single bunch) current to the longitudinal microwave instability was very low. Above threshold, there was small pulse-to-pulse variation in bunch length and longitudinal phase of the extracted beam, that was amplified in the linac, and that made it almost impossible to operate the collider.

The SLC damping ring vacuum chambers seem to have been designed with little concern as to their impedance effects. Or it may have been that, at the time, small vacuum chamber objects and transitions were not considered especially dangerous from an impedance point of view. At about the same time (the early 1980’s) an impedance upgrade to SPEARII involved removing from the ring large objects, such as RF cavities and beam separators, with seemingly little concern for smaller objects. (This left the ring more inductive, resulting in longer bunches, which is probably why mini-beta failed to increase the luminosity [10].) It turned out (as we will show) that for the SLC damping rings it was, indeed, small objects that dominated the original impedance.

The SLC damping ring vacuum chamber has had three incarnations. After beginning operation with the old chamber, when the microwave instability was recognized as limiting performance (in the late 1980’s), the bellows were sleeved, resulting in the old chamber. In 1994 the entire chamber in both rings was removed, and a new, low impedance vacuum chamber was installed (the current or new chamber).

In this report we present, in chronological order for all three versions of vacuum chambers, the original calculations of the impedance and its expected effect on the beam, and the comparison with measurement. At the time the calculations were done we were somewhat limited by computing power and programs for calculating wakefields; much better calculations were possible already by the time of DAΦNE (described in the next section), and even more is possible today. Nevertheless, even with our sometimes crude calculations we were able to reproduce (admittedly, at times in hindsight) many important features of bunch lengthening and the microwave instability found in measurement, and also to provide insight into the physics behind this complicated phenomenon.

The original impedance calculations come from Ref. [11],[12], for the original/old rings, Ref [13] (with C.-K. Ng) for the new ring. The instability calculations can be found in Refs. [14] (old ring) and [15] (new ring). We apologize that many figures are not clear; their originals could not be found, and they were extracted from pdf files of reports. Selected SLC damping ring parameters are given in Table 1.

Table 1: Selected SLC damping ring parameters. Note that, sometime after the new ring chamber was installed, the nominal bunch length and energy spread were increased (by 6%), and the damping time reduced, by a modification of the damping partition numbers [16].

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
<th>Units</th>
</tr>
</thead>
<tbody>
<tr>
<td>Energy</td>
<td>1.15 GeV</td>
<td></td>
</tr>
<tr>
<td>Circumference</td>
<td>35 m</td>
<td></td>
</tr>
<tr>
<td>Typical beam pipe radius</td>
<td>1 cm</td>
<td></td>
</tr>
<tr>
<td>RF frequency</td>
<td>714 MHz</td>
<td></td>
</tr>
<tr>
<td>Nominal RF voltage</td>
<td>0.8 MV</td>
<td></td>
</tr>
<tr>
<td>Nominal rms bunch length</td>
<td>5 mm</td>
<td></td>
</tr>
<tr>
<td>Nominal rms energy spread</td>
<td>0.07 %</td>
<td></td>
</tr>
<tr>
<td>Nominal synchrotron frequency</td>
<td>99 kHz</td>
<td></td>
</tr>
<tr>
<td>Synchrotron radiation damping time</td>
<td>1.7 ms</td>
<td></td>
</tr>
</tbody>
</table>

Original vacuum chamber

The study of the SLC damping ring impedance began with blueprints of the vacuum chamber, to obtain dimensions, which were then used to calculate wakefields. We felt it was better to start from first principles than use simplified models, such as the $Q = 1$ resonator model (see e.g. [17]), as is sometimes done. The goal of the impedance calculations was to obtain two things: (i) an understanding of the relative importance of various objects,
and (ii) a pseudo-Green function wake—that represents, as accurately as possible, the interaction of all objects—that can be used in potential well, threshold, and tracking simulations.

The original ring vacuum chambers were composed of 40, essentially cylindrically symmetric, quad chamber segments (20 each of “QF” or “QD” type), separated by bend chambers with a rather rectangular cross-section (see Fig. 1). The quad segments are sketched in Fig. 2. We see bend-to-quad transitions, cavity-type beam position monitors (bpms), bellows, masks, and flex joints. The ring also has special chambers that include such things as 2 two-cell rf cavities, kickers, septa, y-joints, etc.

Figure 1: Cross-section of a bend chamber. The dashed circle shows the size of a quad chamber beam pipe.

Figure 2: Vertical profile of QF segment (top) and QD segment (bottom). Non-cylindrically symmetric parts are indicated by dashes.

The characterization of a ring impedance as inductive, resistive, or capacitive comes from electrical circuit analogies (an early usage was by Ha"issinski [18]). A ring can be said to be inductive if the induced voltage can be written as \( V_{\text{ind}} \approx -L \frac{dI}{dt} \), with \( I(t) \) the bunch current and \( L \) a constant (the inductance), and it can be characterized as resistive if \( V_{\text{ind}} \approx RI \), with \( R \) a constant (the resistance), etc. An inductive ring means potential well bunch lengthening and increased tune spread with current, a resistive ring has little of either (see, e.g. Refs. [12],[19]). Individual vacuum chamber objects can often also be characterized as inductive, resistive, etc. Small objects or gradual perturbations—bellows, masks, transition—tend to be inductive at normal bunch lengths; larger objects, such as RF cavities and cavity beam position monitors (bpm’s), tend to be resistive or have large resistive components. Note that for most vacuum chamber objects these simple models, at best, only approximately describe their wakefields.

Wakefields of vacuum chamber objects in the SLC damping rings were obtained using early versions of MAFIA2D and MAFIA3D, which are finite difference mesh programs that compute the wakefield of a gaussian bunch in the time domain [20]. For the many inductive objects in the SLC damping ring the wake was calculated for a nominal bunch and then fit for the effective inductance \( L \).

Table 2: Inductive vacuum chamber objects in the original SLC damping rings.

<table>
<thead>
<tr>
<th>Type</th>
<th>Single Element Inductance ( L/\text{nH} )</th>
<th>Contribution in Ring Factor</th>
<th>Number</th>
<th>Contribution in Ring ( L/\text{nH} )</th>
</tr>
</thead>
<tbody>
<tr>
<td>QD bellows(^*)</td>
<td>0.62</td>
<td>1.0</td>
<td>20</td>
<td>12.5</td>
</tr>
<tr>
<td>QD and QF masks</td>
<td>0.47</td>
<td>1.0</td>
<td>20</td>
<td>9.5</td>
</tr>
<tr>
<td>QD &amp; QF trans.</td>
<td>0.52</td>
<td>0.9</td>
<td>20</td>
<td>9.3</td>
</tr>
<tr>
<td>ion pump slots</td>
<td>1.32</td>
<td>0.1</td>
<td>40</td>
<td>5.3</td>
</tr>
<tr>
<td>kicker bellows(^*)</td>
<td>2.03</td>
<td>1.0</td>
<td>2</td>
<td>4.1</td>
</tr>
<tr>
<td>flex joints</td>
<td>0.18</td>
<td>1.0</td>
<td>20</td>
<td>3.6</td>
</tr>
<tr>
<td>1&quot; bpm trans.</td>
<td>0.10</td>
<td>0.8</td>
<td>40</td>
<td>3.3</td>
</tr>
<tr>
<td>other</td>
<td></td>
<td></td>
<td></td>
<td>2.4</td>
</tr>
<tr>
<td>Total</td>
<td></td>
<td></td>
<td></td>
<td>50.0</td>
</tr>
</tbody>
</table>

*Shielded in the late 1980’s by the addition of sleeves.

To obtain the pseudo-Green function wake, calculations were performed using a \( \sigma_z = 1 \) mm gaussian bunch. To properly account for the interaction of neighboring vacuum chamber objects the QF and QD segments—containing the most important impedance objects—were each calculated as one piece. As far as obtaining an accurate Green function we were fortunate that the dominant contributors to the impedance were the bellows, the masks, and the transitions. The total inductance, 50 nH, equivalent to \( |Z/n| = 2.6 \) \( \Omega \), was very large. The rf cavities are resistive, with an effective resistance of 411 \( \Omega \).

Convolving (minus) the Green function with a 6 mm gaussian bunch we obtain the bunch wake shown in Fig. 4 (here a negative value indicates voltage loss). We see that, though largely inductive, the ring wake has a significant resistive component. Performing the Fourier transform of the Green function we obtain the impedance; we find that \( |Z/n| \) rises to a maximum at 15 GHz (above cut-off) due to the bellows, and then drops to become numerical noise beyond 30 GHz (see Fig. 5). We see that, by simply shielding the bellows, the impedance can be significantly reduced (the dots in the figure give the remaining impedance).
To obtain the steady-state bunch shape below threshold we used the Green function and numerically solved the Haïssinski equation [18]. To get the average shape above threshold, we used the same method but first scaled the nominal bunch length parameter $\sigma_{z0}$ by the energy spread increase, as obtained from measurements (described below). The scaling turns out to be $\sim N^{1/3}$.

**Measurement [21]** During SLC operation the beam, once extracted from the ring, passed through the ring-to-linac (RTL) transfer line on its way to the linac. Using a digitized phosphor screen located in a dispersive region of the RTL one could measure the energy spread or—after inducing a longitudinal-energy correlation in the beam—the bunch length of the beam. The current-dependent centroid shift or, equivalently, the parasitic mode losses, were obtained by measuring the RF component of a bunch intensity signal as the beam current was gradually scraped away.

The measured bunch length, energy spread, and centroid shift as functions of bunch population are shown in Fig. 6. The energy spread appears to be independent of current up to $N_{th} = 1.5 \times 10^{10}$, and then increases as $\sigma_E \sim N^{1/3}$. We note that bunch lengthening (frame a) is quite pronounced: the full-width-at-half-maximum (FWHM) length, $z_{fwhm}$ has doubled by $N = 3 \times 10^{10}$. The fact that $z_{fwhm} > 2.355\sigma_z$ ($\sigma_z$ is rms length) indicates that the beam is more bulbous than a gaussian, which is consistent with an inductive impedance [19]. The potential well calculations are given by lines in Figs. 6 (dashed lines above threshold); we see good agreement with measurement. in Fig. 7 we present selected measured bunch shapes (the plotting symbols) and their comparison with calculation (the lines).
Old ring

In the original SLC damping rings, the microwave instability was encountered at $N = 1.5 \times 10^{10}$. Our simulations showed that the bellows were a dominant impedance source, and that the instability threshold could be increased a factor of 2 by shielding them. Sleeves were inserted in the bellows of both rings. This version of the machine we call the old ring.

Shielding the bellows also made the ring somewhat less inductive, reducing the estimated strength of inductors from 50 nH to 33 nH. Again a Green function was generated and potential well calculations were performed; the calculated bunch length with current was not very different from before. In addition, macro-particle tracking of longitudinal phase space was performed, in order to simulate the beam behavior above threshold \([22]-[26]\). In this calculation the longitudinal position and energy of a few hundred thousand macro-particles were tracked. On each turn, each particle’s energy was modified to include effects of the RF wave, Robinson damping, quantum excitation, radiation damping, and the wakefield; and then each particle’s position was modified through the momentum compaction factor.

A tracking example using the old ring Green function is given in Fig. 8. Shown is the turn-by-turn ($N_t$ is turn number) skew moment when $N = 3.5 \times 10^{10}$ (a), and the rms when $N = 5.0 \times 10^{10}$ (b).

The Fourier transform (FT) of the turn-by-turn skew moment of the calculations for the same two currents is given in Fig. 9. At some currents, such as at $N = 3.5 \times 10^{10}$, we obtain an extremely narrow resonance. The position of the peaks of the resonance in the Fourier transform, normalized to the synchrotron tune, as function of $N$ is plotted Fig. 10. Beginning with a calculated threshold $N_{th} = 2 \times 10^{10}$ the resonance frequency grows from $2.5\nu_{s0}$ ($\nu_{s0}$ is nominal synchrotron tune) nearly linearly with slope $0.27\nu_{s0}/10^{10}$.

A tracking example using the old ring Green function is given in Fig. 8. Shown is the turn-by-turn ($N_t$ is turn number) skew moment when $N = 3.5 \times 10^{10}$ (a), and the rms when $N = 5.0 \times 10^{10}$ (b).

The Fourier transform (FT) of the turn-by-turn skew moment of the calculations for the same two currents is given in Fig. 9. At some currents, such as at $N = 3.5 \times 10^{10}$, we obtain an extremely narrow resonance. The position of the peaks of the resonance in the Fourier transform, normalized to the synchrotron tune, as function of $N$ is plotted Fig. 10. Beginning with a calculated threshold $N_{th} = 2 \times 10^{10}$ the resonance frequency grows from $2.5\nu_{s0}$ ($\nu_{s0}$ is nominal synchrotron tune) nearly linearly with slope $0.27\nu_{s0}/10^{10}$.
Figure 10: The positions of the major peaks in the Fourier transform of the skew moment, normalized to the synchrotron tune, vs. $N$.

Figure 11: For old ring, the bunch shape at extrema of the mode oscillation, when $N = 3.5 \times 10^{10}$ (the dashed curves give the average, Haïssinski distribution). The head is to the left.

Vlasov equation calculation Another program useful for understanding a longitudinal, microwave instability was written by one of us (K.O.); it solves the time independent, linearized Vlasov equation including the effects of potential well distortion; we will refer to it here as “the Oide Vlasov solver” [5]. The authors end up with an infinite dimensional, linear matrix eigenvalue problem that is truncated to finite size. For a given current the problem is solved, and the appearance of an eigenmode with complex eigenvalue indicates an unstable mode. The program is used to find the threshold, and also the shape and frequency of the unstable mode. One can also approximately go beyond threshold by assuming the average energy distribution remains gaussian, with the rms energy spread increasing to keep the beam just at threshold.

Using the old SLC damping ring Green function, the first unstable mode the program found was at $N = 1.9 \times 10^{10}$ with a frequency of $2.5\nu_{s0}$, and the shape of the mode was also similar to that found by the macro-particle tracking program. Due to the strongly inductive nature of the ring, and its attendant large tune spread, already at $N = 0.5 \times 10^{10}$ modes with different azimuthal and radial mode numbers overlap in frequency; at threshold the instability cannot be described as the collision of two simple modes (a la mode coupling theory [27]).

Figure 12: For old ring, the shape of the unstable mode from two views at $N = 3.5 \times 10^{10}$.

Measurement [28] After the bellows were sleeved the measured bunch length was similar to that of the original ring, only 10% shorter at $N = 3 \times 10^{10}$. The threshold increased by a factor of two to $N_{th} = 3 \times 10^{10}$. At threshold, on a spectrum analyzer set to a revolution harmonic near 20 GHz, a sideband with a frequency shift of $2.5\nu_{s0}$ was observed (the “sextupole” mode); at higher current the frequency shift increased. A second, weaker sideband was found at about twice the frequency of the first one.

At certain currents above threshold, a repeating sequence of growth and relaxation was observed on beam phase and bunch length signals of the stored beam. The growth time was on the order of the synchrotron period, and the relaxation time on the order of the damping time (the “saw-tooth”). A parameter study over RF voltage and beam current found what might be called phase transitions, with different regions of parameter space displaying qualitatively different types of behavior (the “Nose Plot” [29],[30]). For example, under some conditions the saw-tooth disappeared completely, to be replaced by even oscillations.

Comparing with calculations, the measured bunch length with current was still in good agreement. The measured threshold had increased by a factor of two (as in calculation), though the absolute value was still 50% larger than calculated. The slope of the “sextupole” mode frequency with current was in good agreement [31]; the second sideband, however, was not seen in simulation. Also, although continuous oscillation and sawtooth above threshold could both be found in tracking simulations, we would not say that the calculated saw-tooth convincingly corresponded to measurement. Finally, note that the discovery of the saw-tooth instability in the SLC damping rings engendered much theoretical interest; see e.g. Ref. [32].
Current ring

It was practically impossible to operate the SLC beyond threshold $N_{th} = 3 \times 10^{10}$ because of the microwave instability. An additional problem with the impedance was that, at higher currents, the bunch length extended beyond the linear region of the RTL compressor RF wave, and bunch compression became inefficient. So in 1994 a completely new, low impedance vacuum chamber was installed in both rings. The magnets were not changed, so that quad-to-bend transitions were still needed, but the new ones were much smoother (see Fig. 13); the cavity style bpm’s were also not changed. This version of the ring we call here the current or new ring.

Figure 13: The new, smoother quad-to-bend transitions of the current damping ring vacuum chamber. From Ref. [13].

According to our calculations the impedance of the inductive elements for the old ring was 33 nH; with the new chamber inductive elements were eliminated as much as possible, and we estimated the residual inductance at 6 nH [13]. The vacuum chamber had changed character and become resistive. The new Green function wake is shown in Fig. 14. In Fig. 15 we give an example potential well calculation with this wake. We see that the induced voltage is, to good approximation, proportional to the bunch shape. The dashed curve gives the potential well result for a purely resistive impedance with $R = 880 \Omega$, which we see gives almost the same bunch shape.

Figure 14: The Green function wake representing the current SLC damping rings.

Calculations with the new wake predicted a shorter bunch, and according to particle tracking (again with the damping time artificially reduced by a factor of 10) the threshold appeared to move to $N = 5 \times 10^{10}$. When the new machine was turned on, it was found that bunch lengthening was indeed reduced; the threshold, however, also went down (which came as a shock).

Weak instability

About the same time one of us (K.O.) was studying the microwave instability in a ring with an idealized resistive plus inductive impedance [33]; it was found that a machine with a purely resistive impedance is unstable at any current. Unlike F. Sacherer’s mode coupling instability, which can be described as two modes with different azimuthal mode numbers colliding, this new type of instability can be described as the collision of two modes with identical azimuthal mode numbers but different radial mode numbers. A simple, double water bag model that describes this type of instability was also developed [34]. Unlike the normal mode coupling instability, it is a weak instability, in that it can be suppressed by a small amount of tune spread—as would be introduced by adding a small amount of inductance to the ring—through Landau damping. We call this instability weak as opposed to F. Sacherer’s strong instability. K.O.’s study suggested that the instability in the new SLC damping ring might be such a weak instability, and our subsequent simulations supported this idea.

Simulations

Unlike the old, strong SLC damping ring instability, a weak instability is very sensitive to damping time; this parameter, therefore, could not be artificially lowered in simulations of the new ring to save computing time\(^1\). To keep the total running time manageable, the number of macro-particles instead was reduced by a factor

---

\(^1\)It was one of us (K.B.) who had not checked this point in the earlier, new damping ring simulations. The redeeming feature of this mistake is that, had we known beforehand that the threshold would be reduced, we might not have built the new chambers, and consequently not have achieved the later, improved luminosities.
of 10—-to 30,000—as compared to before. Below threshold and to find threshold this worked fine; above threshold, however, the combination of small number of macro-particles, long damping time, and type of impedance resulted in large fluctuations. This can be seen in Fig. 16, where we plot the turn-by-turn rms energy spread just above threshold (a), and at a higher current (b). Even though there are large fluctuations in the moments of the beam distribution above threshold, we took the average result over the last damping time to estimate the expected average property.

Figure 16: The turn-by-turn rms energy spread obtained by tracking, just above threshold (a) and at a higher current (b).

According to macro-particle tracking the threshold \( N_{th} \approx 1.15 \times 10^{10} \), but the result is very sensitive to inductance. By adding a small inductance of 2 nH (\(|Z/n| = 0.1 \Omega \)) to the impedance, the threshold can be raised by \( 1 \times 10^{10} \) (in agreement with Ref. [33]). When artificially changing the damping time \( \tau_d \), we find that \( N_{th} \sim \tau_d^{-1/2} \), which implies a growth \( \sim e^{\alpha N^2 t} \) (\( \alpha \) a constant, \( t \) time), also in agreement with Ref. [33]. The threshold and the average beam properties as functions of current above threshold are in good agreement with the Vlasov equation results. The unstable mode is clearly a quadrupole mode (see Fig. 17). The unstable mode frequency is just under \( 2\nu_s/10 \) at threshold, and then varies with current with a slope of \( \sim 0.06\nu_s/10^{10} \), also in good agreement with calculation.

Figure 17: Unstable mode shape at \( N = 2 \times 10^{10} \) as obtained by the Vlasov method. The bunch head is to the right, higher energy is down.

**Measurement [35]** The impedance effects in the current damping ring have been extensively studied through measurement, resulting in one, and the significant part of another, PhD.: for bunch length measurements, primarily using a streak camera [36], and for a detailed study of the properties of the unstable mode above threshold [16]. Bunch length measurements clearly showed that bunch lengthening had been reduced by the introduction of the new chambers (see Fig. 18). The threshold was found at \( N_{th} = 1.5-2 \times 10^{10} \). Bunch length and synchronous phase were in good agreement with calculations (when 2 nH inductance was added to our Green function); the bunch shape was found to be consistent with a resistive impedance, and in agreement with calculation. The unstable mode, at threshold, had a frequency of \( 1.77\nu_s/10 \), and varied with current with a slope of \( -0.06\nu_s/10^{10} \), also in good agreement with calculation.

Measurements of time dependent (saw-tooth) behavior for different currents are shown in Fig. 19. The instability involved the movement from the equilibrium shape of only a few percent of beam particles [16]. It appears that the amplitude of instability in the new ring was less than in the old (old ring saw-tooth measurements were not calibrated), since with the old ring we were limited to threshold, but with the new ring we ran routinely at more than twice threshold, at \( N_{th} = 4.5 \times 10^{10} \).

**Discussion**

Recent programs for calculating impedances/wakefields and for tracking longitudinal phase space in storage rings are much improved. We would briefly like to mention one: a numerical method for solving the Vlasov-Fokker-Planck equation for longitudinal phase space in storage rings has been developed by Ellis and Warnock [8]. The program seems better able to avoid fluctuations in beam properties above threshold that we found, e.g., with macro-particles in Fig. 16b. Their program was applied to our Green function wakes for the old and current SLC damping ring vac-

---

Figure 18: Bunch length (FWHM/2.35) vs. current measurements, comparing the new with the original (here called “old”) chambers. The new measurements were performed using a streak camera. From Ref. [36].

Figure 19: Unstable mode frequency vs. current measurements for different currents in the SLC damping ring.
relations, and the period of bursting envelope are in good agreement with measurement for the new ring, other properties are not (e.g. the saw-tooth behavior does not disappear again at higher currents, as seen in Fig. 19). An example with agreement: in Fig. 20 we see a simulated oscilloscope trace, obtained by the program, which is meant to be compared to the third curve from the top in Fig. 19.

So finally, how can we understand the reduction of the measured threshold when the SLC damping ring impedance was reduced? In the old, inductive machine there was a strong instability observed at \( N_{th} = 3 \times 10^{10} \), and we expect to have increased this threshold when the impedance was reduced. However, in an inductive machine there is a large incoherent tune spread that will Landau damp weak instabilities which might otherwise appear at lower currents. By removing mostly inductive elements, and thereby changing the character of the ring to a resistive one, we have removed this tune spread, and presumably are now able to observe one of these weaker instabilities.

If one could reduce a ring impedance by a scale factor, \( \alpha \), one would then have confidence that the dependence of wakefield effects on current shifts up by \( 1/\alpha \). But such a change in ring impedance is generally not realistic. The Boussard criterion suggests that the threshold to the strong instability depends on \( |Z/\alpha| \), but it is known that this is a simplification, that the character of the impedance is also important; and the Boussard criterion says nothing about the weak instability. When reducing an impedance to ameliorate wakefield effects one generally needs to do more careful analysis. For the current SLC damping rings how could we raise the instability threshold? We believe that the instability in the new rings is very sensitive to a small amount of Landau damping. It could be damped by adding a weak, higher harmonic cavity, or by reinserting a small amount of inductance, by e.g. introducing a short bellows, or beam pipe with many small holes. However, one must take care not to overdo it, thereby bringing back down the threshold to the strong instability.

**Conclusion**

We have reviewed our development in understanding of the longitudinal impedance and microwave instability in the SLC damping rings, from the original, to the old, and finally to the current (or new) versions of the ring vacuum chambers. Our calculations were somewhat crude, when compared to what can be done today. Nevertheless, the calculations were a useful complement to measurement, giving reasonable agreement with measurement and insight into the instability, whether in the old, inductive, or new, resistive rings. In the process a new kind of instability—the weak instability—was discovered; it can now be considered to be reasonably well understood.

**DAΦNE**

**Vacuum chamber RF design**

The vacuum chamber RF design of the Frascati e⁺e⁻ \( \Phi \)-factory DAΦNE[38] is a good example of successful benchmarking of impedance codes.

First, the DAΦNE vacuum chamber is complicated. Despite a short collider circumference of about 97 m each ring accommodates all components typical for a multibunch high current collider. The rings contain:

1. two common 10 m long Y-shape interaction regions;
2. four 5 m long narrow gap wigglar vacuum chambers (the wigglers are used for the emittance control and enhancement of the radiation damping);
3. straight sections for allocation of RF cavities, injection kickers, longitudinal feedback kickers, transverse feedback kickers etc;
4. tapers connecting straight sections, bending arcs, wigglar sections, interaction regions;
5. many other components as bellows, flanges, valves, vacuum ports etc.

Second, due to high circulating currents the vacuum chamber was designed to avoid beam instabilities and excessive power losses. New designs and novel ideas were adopted for almost all principal vacuum chamber components: RF cavities [39, 40], shielded bellows [41], longitudinal feedback kickers [42], BPMs [43], DC current monitors [44], injection kickers [45], transverse feedback kickers and others [46]. For example, longitudinal feedback kickers similar to the DAΦNE kicker are routinely used in more than 10 operating colliders and synchrotron radiation sources.

At present, the design single bunch current of 44 mA has been largely exceeded. About 200 mA were stored in a single bunch, while in the multibunch regime 2.4 A of stable beam current were accumulated in the electron ring and about 1.3 A in the positron one.

The following impedance codes were used for impedance and wake field calculations in DAΦNE: ABCI [47], URMEI [48] and OSCAR2D [49] were used in simulations of azimuthally symmetric structures, while the impedance of 3D objects was calculated by MAFIA [20] and HFSS [50]. As a complement to MAFIA and HFSS, the POPBCI code [51] was used to characterize the higher order mode content in the DAΦNE (rectangular waveguide loaded) RF cavity [52].

Impedance bench measurements were carried out for almost all critical vacuum chamber components. Generally, an agreement between the measurement results and the impedance code simulations is satisfactory. Examples of such a comparison are given in a review paper [53], and more details can be found in Refs. [39]-[46].

The total collider impedance estimate obtained with the impedance codes very reliably predicts such important aspects of beam dynamics in DAΦNE as: bunch shortening, bunch shape and the threshold to the microwave instability.

**Bunch lengthening**

Bunch lengthening simulations for DAΦNE were performed much before the collider commissioning. The overall short range wake used in the numerical tracking was calculated by adding up contributions of almost all vacuum chamber discontinuities that were estimated analytically or numerically assuming a 2.5 mm gaussian distribution [54], see Fig. 21.

The tracking method is essentially the same as that successfully used in the bunch lengthening simulations for the SLC damping rings [14], SPEAR [25], PETRA and LEP [24]. It consists in tracking the motion of $N$ superparticles in longitudinal phase space over 4 damping times. The turn-by-turn equation of macroparticle motion includes lattice dispersion, radiation damping, stochastic quantum excitation, interactions with the RF field and the wake of all leading macroparticles.

A comparison between bunch lengthening simulations and bunch length measurements using two different methods is shown in Fig. 22. In the beginning of DAΦNE commissioning a signal from a broadband button was used for bunch length measurements (blue points in Fig. 22). The resulting bunch distribution was found by processing the signal picked up by the button, taking into account the button transfer impedance and the attenuation of the cables connecting the button to a sampling oscilloscope [55]. An example comparison between the simulated bunch shape and the processed signal is shown in Fig. 23.

**Figure 21:** Wake potential of a 2.5 mm long gaussian bunch, which is used in bunch lengthening simulations for DAΦNE.

**Figure 22:** Comparison of bunch lengthening simulations (green line) with bunch length measurements performed with a BPM (blue circles) and a streak camera (red squares).

Later installation of a streak camera made it possible to measure the bunch length in both the electron and positron rings simultaneously [56] (red squares). The bunch profiles at different currents as acquired by the streak camera are shown in Fig. 24. The difference in bunch shape in the two rings is due to an additional inductive impedance in the electron ring: the ion clearing electrodes [57].

The calculated DAΦNE impedance is suitable for predicting bunch behavior not only in the lengthening regime, but also when the short-range wake becomes focusing (bunch shortening). Recently, it has been proposed to use a lattice with a negative momentum compaction factor to
increase the luminosity in DAΦNE [58]. For such a lattice the short range wakes are focusing and the bunch shortens until the microwave threshold is reached. An experimental lattice with a negative momentum compaction factor has already been tried in both DAΦNE rings [59]. Fig. 25 shows bunch length as function of current and Fig. 26 gives typical bunch distributions in the positron ring with negative momentum compaction. Again, we can find an agreement between the measurements and the predictions of [58] as far as the bunch length and shape are concerned.

Figure 23: Comparison of the (processed) BPM signal (dotted line) with simulated bunch profile (solid line). The bunch head is to the left.

Figure 24: Typical measured bunch distributions in the positron (left) and electron (right) rings. The head is to the left.

Figure 25: Measured bunch length in the DAΦNE positron ring as function of bunch current for positive (red squares) and negative (blue circles) momentum compaction factor.

Figure 26: Typical measured bunch distributions in the DAΦNE positron ring with a negative momentum compaction lattice. The head is to the left.

**Microwave instability**

The coupling between longitudinal coherent modes in a bunch is the driving source of the microwave instability. Different azimuthal modes may couple if their natural frequencies are shifted by amounts comparable to the synchrotron frequency ("strong" instability) while radial modes having the same azimuthal number can couple already for much smaller frequency shifts ("weak" instability). The coherent shift is due to the interaction between the bunch and the machine impedance.

In order to study the microwave instability for DAΦNE again we assumed the machine wake function calculated by the impedance codes shown in Fig. 21. The simple analytical model that we followed treats mode coupling as a splitting of each azimuthal mode in two radial modes. The model is based on approximating the real bunch distribution by a double water bag distribution [34]. By substituting this distribution into the Vlasov equation we solved the resulting eigenvalue system after truncating it (keeping only first 9 azimuthal modes). According to this study, the coupling of radial modes with low azimuthal mode number drives the microwave instability at DAΦNE. For instance, as is seen in Fig. 27, at an RF voltage of 100 kV the lowest thresholds are given by the coupling of radial quadrupole modes at a bunch current of 24 mA and sextupole modes at 28 mA.

Figure 27: Frequencies of radial bunch modes with lowest azimuthal mode numbers \( m = 1, 2, 3 \) as a function of bunch current.
Experimentally, at a voltage of 100 kV and with the same momentum compaction factor as considered in the analytical model, the quadrupole mode was clearly observed with a spectrum analyzer (see Fig. 28). In the positron ring the threshold current at which the mode signal first appears is about 26 mA, a value that is surprisingly close to the prediction of the simplified analytical model.

In practice, in the positron ring it was possible to push the threshold to higher currents by varying the RF voltage and increasing the momentum compaction factor. In the electron ring the quadrupole mode instability problem was more severe, presumably due to the higher broadband coupling impedance. The instability had been limiting the maximum stored current in the e⁻ ring for a long time, leading to injection saturation, background problems, beam-beam blow up and lifetime reduction. The voltage variation and momentum compaction increase gave only a small increase in the maximum storable current. The problem has been finally solved by tuning the DAΦNE longitudinal feedback system (that initially meant to damp only dipole oscillations) in such a way as to give different longitudinal kicks to the head and tail of bunches [60].

**DAΦNE Accumulator Ring**

Yet another example of successful impedance code application is the DAΦNE Accumulator Ring, a small booster ring in the DAΦNE injection chain [61]. As in the case of the main rings, the coupling impedance was estimated well in advance of measurements. The longitudinal coupling impedance of 3.5 Ω inferred from the bunch lengthening measurements is in a good agreement with numerical impedance simulations [62]. Fig. 29 shows the calculated wake potential for a 5 mm long gaussian bunch, and its Fourier transform (the impedance). Fig. 30 gives a comparison between calculated and measured bunch lengths at different RF voltages and bunch currents.

The measurement of the shift of the transverse betatron tunes and their synchrotron sidebands versus bunch current has shown that the accumulator ring broadband impedance can be approximated by a $Q = 1$ broad-band resonator model, with a shunt impedance of 70 kΩ/m [63]. This agrees well with earlier analytical estimates and numerical calculations [64].

**APPENDIX: PANEL COMMENTS**

**Conclusions on impedance codes-MZ**

In our opinion, there is no urgent need to develop entirely new software packages for impedance calculations than available today. It would be sufficient to extend already existing tools. The present general-purpose impedance codes, like MAFIA, HFSS, GdfidL and others, have proven their reliability in RF designing and vacuum chamber’s impedance optimization. In our paper we have given only a few examples of this.

It does not mean that the existing codes can solve all...
possible problems arising while calculating the impedance. However, in most cases numerical simulations can be always crosschecked (or substituted) with analytical evaluations and/or experimental measurements results. For particular cases analytical estimates can be helpful in simplifying models to give a possibility of numerical simulations with the existing codes.

Among the practical suggestions for further development of the existing codes we can list a few:

1. possibility to include resistive walls in simulations to take into account:
   - crosstalk between resistive wall and geometric wake fields. Recently this task is getting more important for design of collimation systems for future accelerators, first of all, LHC and linear colliders. - quadrupolar resistive wall wakes responsible of the betatron tune shifts in asymmetric vacuum chambers (measured, for example, in PEP-II, DAΦNE,...)
2. simulations of vacuum chamber elements with thin resistive layers. There are many examples of such components: CTF3 BPMs, DAΦNE ion clearing electrodes, different kinds of kickers etc.
3. further impedance benchmarking is necessary for simulation of components containing frequency dependent materials and/or frequency dependent external loadings
4. possibility of direct impedance calculations at a given frequency.

Comments-KB

Wake Calculation

- As rings become cleaner 3D objects become more important; for short bunch, wakes of long 3D objects are difficult to calculate accurately.
- Short gaussian bunch (used for Green function) filters out high frequencies; what if instability is driven at very high frequencies (e.g. CSR, microbunching instability)?
- For short bunch (microbunch), interaction can occur over long distances (catch-up problem).
- Short bunch, long structure, small features: difficult to calculate accurately (mesh dispersion).

Vlasov Equation Programs

- Oide turned the linearized Vlasov equation into a linear eigenvalue problem; many superfluous stable modes (artifact). Program does not always work.

=> R. Warnock, M. Venturini, G. Stupakov, turned the linearized Vlasov equation into a nonlinear eigenvalue problem; no superfluous modes. More likely to work for difficult wakes, though does not always work.

Tracking

- Tracking above threshold can yield large fluctuations.

=> R. Warnock and J. Ellison have developed a program that solves the Vlasov-Fokker-Planck (VFP) equation; more accurate than simple tracking; can e.g. solve CSR-driven, microbunching instability; can it be benchmarked with measurements?

ACKNOWLEDGMENTS

We acknowledge the debt of many people who were involved in the SLC and DAΦNE—physicists, engineers, operators—that made it possible to come to some understanding of the ring impedances. We thank also A. Chao, for reading the manuscript and giving useful comments and suggestions; F.-J. Decker, for recent discussion about SLC measurements; R. Warnock and B. Podobedov for supplying figures.

REFERENCES


[50] “HFSS, the High Frequency Structure Simulator HP85180ATM.”


Beam-beam Simulations of Hadron Colliders
Tanaji Sen
Fermilab, PO Box 500, Batavia, IL 60510

Abstract
Simulations of beam-beam phenomena in the Tevatron and RHIC as well for the LHC are reviewed. The emphasis is on simulations that can be closely connected to observations.

I. INTRODUCTION
Beam-beam phenomena have been studied intensively ever since the first storage ring collider started operation. These interactions have limited the achievable luminosity and beam intensities in both hadron and lepton colliders. However the nature of the limitations in these two classes of colliders is sufficiently different that distinctly different analytic and simulation tools have been developed. In this review I will focus on recent developments in beam-beam simulations for hadron colliders. The emphasis will be on connecting simulation results to observations; almost nothing will be said about the simulation techniques.

II. USES OF SIMULATIONS
Simulations of beam dynamics can be useful for different reasons. At the design stage of an accelerator they can be useful in guiding the proper choice of beam and machine parameters. Examples would be the several simulations done for the LHC. Historically this has been the most important contribution of simulations. Simulations can also be useful in understanding beam observations in existing accelerators. I will present several examples for the Tevatron later in this report. Finally, simulations could be used to improve the performance of an existing accelerator. This is perhaps the most demanding requirement. To date there are few examples of beam-beam simulations of hadron colliders having been useful in this respect. This situation is changing due to advances in physics modeling and computing power. The easier availability of parallel computing makes it possible to include more details of an accelerator in the simulation model. It is also easier to calculate those beam parameters that are routinely measured in every store but are computationally demanding, such as emittances and lifetimes.

III. HADRON COLLIDERS
I will consider beam-beam phenomena in the two existing hadron colliders Tevatron and RHIC as well as the future LHC. HERA is a hybrid (a lepton-hadron collider) and interesting beam-beam observations have been reported recently [1]. However the phenomena appear to be closely connected to the dynamics of the lepton beam.

In the Tevatron, protons and anti-protons circulate within the same beam pipe and collide at two interaction points (IPs) B0 and D0. The ratio of anti-proton intensities to proton intensities is about 1:7 at present, thus the beam-beam phenomena are in the classical “weak-strong” regime. Nonetheless we observe that anti-protons do influence the protons. At injection energy, both beams experience only long-range interactions while at collision there are two head-on collisions and seventy long-range interactions per bunch. These interactions limit beam lifetimes at all stages of the Tevatron’s operational cycle. RHIC collides polarized protons and at other times, heavy ions such as gold. Both beams have similar intensities and the beam-beam phenomena are in the classical “weak-strong” regime. There are two rings and typically each bunch experiences 4 head-on collisions and 2 long-range interactions. Coherent modes excited by these collisions have been observed in RHIC. The LHC will collide proton beams of equal intensities at two high luminosity IPs and also at two low luminosity IPs. Beam-beam phenomena there will be in the “strong-strong” regime. In addition, each bunch experiences several long-range interactions around each IP. These interactions are expected to hurt beam quality and several simulations have addressed their expected impact. Compensation of these long-range beam-beam interactions with wires
has been proposed for the LHC and initial tests with wires acting on a single beam in the SPS have been carried out. Plans are now underway to test this principle in an existing collider with two beams.

IV. SIMULATION CODES

Until recently most hadron collider beam-beam simulation codes were used to calculate tune shifts with amplitude and dynamic apertures (DAs). These quantities are only measured in dedicated beam studies. In recent years several weak-strong simulation codes have been written to calculate quantities that are routinely observed such as emittances and lifetimes. Strong-strong codes typically calculate the mode spectrum and also emittance changes. Table I lists some of the simulation codes and their purpose.

The codes used at FNAL for lifetime calculations are all parallelized codes. DA codes are typically single processor codes but can be “trivially” parallelized. The references should be consulted for more details on these codes.

<table>
<thead>
<tr>
<th>Simulation Code</th>
<th>Purpose</th>
</tr>
</thead>
<tbody>
<tr>
<td>FNAL</td>
<td>Lifetime calculations</td>
</tr>
<tr>
<td>BBSIM [2], LIFETRAC [3], PLIBB [4], BEAMBEAM3D [5]</td>
<td>Lifetime calculations</td>
</tr>
<tr>
<td>MAD [6], SIXTRACK [7], TEVLAT [8]</td>
<td>Dynamic aperture</td>
</tr>
<tr>
<td>RHIC</td>
<td>Coherent modes</td>
</tr>
<tr>
<td>BBDEMO2C [9], BEAMBEAM3D</td>
<td>Coherent modes</td>
</tr>
<tr>
<td>LHC</td>
<td>Diffusion, Dynamic aperture</td>
</tr>
<tr>
<td>WSDIFF [10], SIXTRACK, MAD</td>
<td>Diffusion, Dynamic aperture</td>
</tr>
<tr>
<td>BEAMX [11], COMBI, BEAMBEAM3D</td>
<td>Coherent modes</td>
</tr>
</tbody>
</table>

Table 1: Beam-beam simulation codes. The numbers in [] are the reference numbers.

V. TEVATRON: OBSERVATIONS & SIMULATIONS

Beam-beam interactions have limited the performance of the Tevatron since the start of Run II in 2001. Several improvements have significantly reduced the beam loss due to these interactions. At present the performance limitations are not severe but there are observable influences due to these interactions. See reference [12] for a discussion of beam-beam phenomena in the Tevatron. Here I will choose some observations that have been qualitatively reproduced in simulations.

In early 2003 the Tevatron was operated without transverse dampers but at high chromaticity in order to keep the beams stable. At injection energy, typical chromaticities were about 8 units in both planes. During this period we observed a strong reduction in the anti-proton emittances soon after they were injected.

Figure 1 shows the emittance of 4 anti-proton bunches measured with flying wires 10 times at injection energy during a store in April 2003. These 4 bunches were injected first among the 36 bunches enabling several measurements of their emittance. The emittance of all bunches fell sharply right after injection before reaching an asymptotic value during this time of 15 minutes that the beams were at injection energy. The same phenomenon was observed in several stores in April 2003. This occurred because the DA was much smaller than the initial beam width. Beam outside the DA was lost; the emittance fell until the beam was small enough to fill the DA.

Figure 1: Emittances of anti-protons after injection (April 2003)
The DA extracted from these measurements over several stores for the 4 bunches is shown in Figure 2. The measured DA varied from 3.5-4 $\sigma$ for these bunches. The error bars represent the variations over the measurements. The dynamic aperture was calculated with the programs Sixtrack and Mad. The simulation model included the known multipole moments in all the arc magnets but not the alignment errors. Chromaticities were set to 8 units in both planes. The expected differences between the model optics and the real machine optics are about 20%. Figure 3 shows the calculated DAs as a function of proton intensity. Despite the differences between the model and the machine optics, the DAs calculated by both codes were around 4 $\sigma$ at the proton intensities in the machine – close to the measured values. The impact of the nonlinearities was strong enough that the inaccuracies of the linear optics were less important. The agreement between calculations and measurements was not as good at smaller values of the chromaticity. The calculated DAs fell relatively slowly with chromaticity while measurements at chromaticities of (4,2) units (after dampers were operational) showed no reduction in the initial emittance implying that the DA was significantly greater than the values shown in Figure 2.

Several codes were used to calculate anti-proton lifetimes at injection where the measured lifetime is a few hours. The statistical accuracy from a numerical estimation of a 1hr lifetime in the Tevatron is $\sim$10% if more than $2\times10^{10}$ particle-turns are used. If all the multipole moments in the magnets are included in a simulation model, the computational time required for tracking these many particle-turns is prohibitively large even with the use of multiple processors. Lifetime simulations at injection have therefore included only the beam-beam nonlinearities.

The dependence of the lifetime on the vertical chromaticity was calculated with the code PlibB [4] since observations showed that the lifetime was very sensitive to this parameter. Figure 4 shows the calculated lifetime for anti-proton bunch 1 at injection as a function of the vertical aperture for
different values of the vertical chromaticity. The horizontal chromaticity was fixed at 2 units. At vertical chromaticities \( \geq 4 \) units, the lifetimes are in the range of 1-2.5 hrs and are not very sensitive to the aperture. But at chromaticities \( \leq 2 \) units, the lifetime increases sharply as the aperture increases from 3 to 6 \( \sigma \) before leveling off. The physical aperture at injection is \(~6\sigma~\) - it is interesting that the simulation predicts a significant jump in lifetime as the chromaticity is lowered from 4 to 2 units. Once the transverse dampers were commissioned in the Tevatron, the chromaticities were lowered to \((4,2)\) units and the anti-proton lifetime improved to an average of around 5 hrs at injection.

After both beams are loaded, they are accelerated to 980 GeV in about 84 seconds. There are beam losses during this stage, typically less than 10%. Dedicated machine studies have shown that proton losses during acceleration are not affected by anti-protons while the anti-proton losses increase in the presence of the protons. So far no simulations have been performed to follow the acceleration process, mainly because the losses are typically small.

After reaching flat top, the optics is changed to reduce the \( \beta^* \) at B0 and D0 to 0.35m. In the early stages of Run II there were large anti-proton losses between 10-25% during the portion of the beta squeeze when the helices changed polarities. The minimum separation had dropped to less than 2 \( \sigma \) at this point. When the separator voltages were changed to increase the minimum beam separation to more than 3.5 \( \sigma \), anti-proton losses also dropped significantly. Now beam losses during the squeeze are typically no more than 2%.

After the beams are brought to collision, the beams experience head-on interactions and long-range interactions. Observations show that both protons and anti-protons are influenced by the other beam. The long store times during collision makes possible detailed beam observations that cannot be made during injection. At 980 GeV, there is enough synchrotron radiation light to image the two beams in a synchrotron light monitor.

![Figure 5: Bunch by bunch orbits (top: hor., bottom: vert.) of anti-protons observed during a store at the synchrotron light monitor.](image1)

![Figure 6: Calculated bunch by bunch orbits of anti-protons](image2)

Figures 5 and 6 show the observed and calculated bunch-by-bunch orbits of anti-protons at the synchrotron light monitor. The calculated orbits reproduce the different patterns of the horizontal and vertical orbits as well as the scale of the orbit shifts.
The commissioning of a high frequency 1.7 GHz Schottky monitor has enabled the measurement of bunch-by-bunch tunes during stores – see Reference [13]. Signals are gated in order to acquire the data from individual bunches and so far the bunch by bunch tunes of either anti-protons or protons have been measured in a single store.

The pattern of measured tunes, a representative sample for a particular store is seen in Figure 7, is reproduced in different stores. The calculated anti-proton tunes, see Reference [14], agree reasonably well with the measured values.

Figure 7: Measured anti-proton tunes bunch by bunch.

An example of this scalloped profile can be seen in Figure 9. The rapid emittance growth of most bunches results in a significant luminosity loss. Empirically the problem has been solved by small changes in the tunes. The Tevatron electron lens was used successfully on one occasion to reduce the emittance growth of a selected bunch [15].

Simulations have attempted to understand the sensitive dependence of the emittance growth on the tunes. An example using the code LIFETRAC is seen in Figure 10. Lowering the horizontal tune by 0.005 reduces the emittance growth of most bunches – roughly in accord with observations. – while the emittance of the last bunch grows more rapidly. The scalloped profile is not seen in the profile at the higher tune but nevertheless the simulation does demonstrate the tune dependence of emittance growth.

Beam lifetimes also show strong variations from bunch to bunch. These lifetimes depend on the beam parameters such as the intensities and emittances. Proton bunch intensities and emittances vary by about 10% across all 36 bunches.

Figure 8: Measured proton tunes bunch by bunch in a store.

Figure 9: Vertical emittance profile of anti-proton bunches soon after the start of a store.

Figure 10: Simulation of emittance growth with tune changes.
Figure 10: Horizontal emittance growth of anti-proton bunches calculated using LIFETRAC. Top: tunes are (0.585, 0.575); Bottom: tunes are (0.580, 0.575) - courtesy of A. Valishev.

As an illustration we consider beam parameters during a store on July 16, 2004 when the highest luminosity of $1.3 \times 10^{32}$ cm$^{-2}$ sec$^{-1}$ was recorded in the Tevatron. Figure 11 shows the proton intensities and emittances over all bunches – the largest variations are in the first 3 or 4 bunches at the head of each train. There was a 10% variation in proton intensities with an average intensity $\sim 250 \times 10^9$. The average horizontal and vertical emittances were 20 and 13 $\pi$mm-mrad respectively with a similar 10% variation in each plane over all bunches.

The variation in anti-proton parameters is considerably greater. Figure 12 shows that the anti-proton bunch intensities and emittances varied by a factor of 2 or more in the same store.

The ratio of proton to anti-proton bunch intensities varied in the range from 4 to 10. Anti-proton bunch emittances are largely determined by the length of time spent circulating in the Accumulator where they are stochastically cooled. These bunches are injected four at a time into the Tevatron. Consequently the last four bunches to be injected typically have the smallest emittances.

These large variations in anti-proton bunch parameters naturally lead to large variations in the bunch luminosities, given approximately by

$$L = \frac{f_{\text{rev}} N_p N_A}{4 \pi \beta^* \varepsilon} H \left( \frac{\beta^*}{\sigma_s} \right)$$

where $f_{\text{rev}}$ is the revolution frequency, $N_p$ and $N_A$ are the proton and antiproton bunch intensities of the colliding bunches and $H$ is the hourglass factor.
Beam lifetimes can be usefully split into two contributions: the dominant one from luminosity and the other from dynamic processes unrelated to luminosity such as beam-beam effects, intra-beam scattering, gas scattering etc. To isolate the lifetime due to dynamic processes, we remove the contribution of the luminosity lifetime as
\[
\frac{1}{\tau_{\text{Dyn}}} = \frac{1}{\tau_{\text{Beam}}} - \frac{1}{\tau_{\text{Lum}}}
\]
and the luminosity lifetime for an antiproton bunch for example is
\[
\frac{1}{\tau_{\text{Lum},A}} = \frac{L\Sigma_{pA}}{N_A}
\]

Here \(\Sigma_{pA}\) is the inelastic proton-antiproton scattering cross-section, assumed to be 70 mbarns at 980 GeV. The luminosity lifetime of an antiproton bunch is independent of the intensity of that bunch.

The bunch-by-bunch dynamic lifetimes \(\tau_{\text{Dyn}}\) are shown in Figure 13 for the July 16, 2004 store and for another store on August 18, 2004. The bunch intensities and emittances for the later store can be seen in Reference [12].
As a first attempt we have calculated the lifetimes of anti-protons assuming design values of the beam parameters and without including the variations in bunch parameters. The multi-particle code BBSIM runs on parallel processors and includes the beam-beam interactions, and random dipole noise to mimic gas scattering. The simulation model in the results presented here assumed linear transport between the beam-beam interactions. More recently chromaticity sextupoles have been added to the model. Lifetimes were estimated by tracking $2 \times 10^4$ particles through $10^6$ turns and the bunch intensities were fitted to an exponential decay curve. Particles whose amplitudes exceed specified apertures are flagged so it is possible to estimate lifetimes as a function of the aperture.

Figure 14: Lifetime of anti-proton bunch 1 calculated with BBSIM at collision as a function of the physical aperture.

The dependence of lifetime on the aperture is shown in Figure 14 for antiproton bunch 1. Figure 15 shows the relative lifetimes at $8 \sigma$ for 12 bunches normalized to the lifetime of bunch 1. The results using this model predict that the lifetime of anti-proton bunch 1 is significantly greater than that of the other bunches with only small variations between the other bunches. This lifetime pattern does not resemble either of the patterns seen in Figure 13. The discrepancies are due to several possible factors - the differences in bunch parameters need to be included, there is likely “cross-talk” between machine nonlinearities and beam-beam nonlinearities and the differences between the linear machine optics and the design optics used in the simulations.

The dynamic lifetimes of protons are shown in Figure 16 for the same two stores as in Figure 13. Again 3-fold symmetry is evident in the latter store but not in the earlier store. These dynamic lifetimes

![Figure 14](image1.png)

![Figure 15](image2.png)

![Figure 16](image3.png)
also vary by roughly a factor of 10 between bunches. Intra-beam scattering has some influence on the dynamic lifetimes but cannot account for the observed variations given that proton bunch intensities and emittances are fairly uniform. We observe that beam-beam effects also have a strong influence on the protons. These losses are attributed mainly to the head-on interactions – after beams are brought to collision, proton losses are higher than anti-proton losses. Analysis shows that the non-luminous losses (losses not related to luminosity) of a proton bunch are greater if the vertical emittance of the colliding anti-proton bunches are smaller – seen in Figure 17.

Figure 17: Correlation of proton non-luminous losses with the vertical emittance of the colliding anti-proton bunches (courtesy – P. Lebrun).

This phenomenon is qualitatively understood and has been observed previously at the SPS [16] and HERA [17]. Protons colliding with a smaller emittance anti-proton bunch experience the peak of the beam-beam force closer to the core of the proton bunch leading to larger losses. A quantitative explanation of this phenomenon is still lacking.

We end this section by a comment on longitudinal diffusion in the two beams. At the start of collisions, the longitudinal emittance of both beams is limited to under 4 eV-sec while the bucket area is 11 eV-sec. Intra-beam scattering leads to diffusion into larger portions of longitudinal phase space. Analysis by A. Tollestrup shows that while protons gradually fill the bucket, anti-protons stay confined to their initial area. Figure 18 shows the time evolution of the longitudinal density in a store as a function of the longitudinal action.

Figure 18: Evolution of the longitudinal density as a function of the longitudinal action. Top: protons, Bottom: anti-protons (courtesy: A.Tollestrup)

This suggests that the beam-beam interactions restrict the longitudinal DA of anti-protons to ~4 eV-sec but there is not a similar limit on the protons. This may be the first such observation of beam-beam imposed limitations on longitudinal dynamics.

Table 2 summarizes the recent beam-beam observations at Fermilab and the corresponding simulation activity. While this is a start, much more needs to be done.
Table 2: List of some important beam-beam observations, whether accompanied by simulations and comment on the simulations.

<table>
<thead>
<tr>
<th>Observations</th>
<th>Simulation</th>
<th>Comment</th>
</tr>
</thead>
<tbody>
<tr>
<td>150 GeV</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Antiproton dynamic aperture</td>
<td>Yes</td>
<td>Agree at high chromaticity</td>
</tr>
<tr>
<td>Anti-proton lifetime</td>
<td>Yes</td>
<td>Depends on chromaticity</td>
</tr>
<tr>
<td>Ramp &amp; Squeeze</td>
<td>No</td>
<td>Simulations only at fixed energy</td>
</tr>
<tr>
<td>Losses during ramp</td>
<td>No</td>
<td></td>
</tr>
<tr>
<td>Losses during squeeze (2001-2002)</td>
<td></td>
<td></td>
</tr>
<tr>
<td>980 GeV</td>
<td></td>
<td></td>
</tr>
<tr>
<td>Impact of larger helices</td>
<td>Yes</td>
<td>Dynamic aperture only</td>
</tr>
<tr>
<td>Bunch by bunch tunes, orbits</td>
<td>Yes</td>
<td>Qualitatively good</td>
</tr>
<tr>
<td>Lifetimes, emittance growth</td>
<td>Yes</td>
<td>In progress</td>
</tr>
<tr>
<td>Proton losses at start of stores</td>
<td>No</td>
<td>No simulations of protons yet</td>
</tr>
<tr>
<td>Longitudinal dynamic aperture of anti-protons</td>
<td>No</td>
<td>Occurs on a long time scale</td>
</tr>
<tr>
<td>Tune scan of lifetimes</td>
<td>Yes</td>
<td>Limited tune scan of dynamic aperture</td>
</tr>
<tr>
<td>Cogging and IP scans</td>
<td>No</td>
<td>Limited benefit</td>
</tr>
</tbody>
</table>

VI. RHIC

RHIC is a two-ring collider that has collided several species of ions. Data has been taken with collisions of gold on gold, deuterons on gold as well as protons on protons. Typically each bunch experiences 4 head-on collisions and 2 long-range interactions per turn. Beam-beam effects are observed to impact the emittance growth and lifetime. Simulations showed that the DA improves close to the SPS tunes and operation with polarized protons indeed showed better beam lifetime at these tunes [18]. At usual proton intensities, the beam-beam parameter is 0.004/IP although in some experiments with 2 head-on collisions, beam-beam parameters of 0.007/IP have been achieved.

Coherent modes driven by the beam-beam interactions have been observed when colliding proton beams of equal intensity. In a dedicated experiment with beams in collision only at a single IP and separated everywhere else, the mode frequencies were measured.

Figure 19: Measured and calculated coherent modes of protons in RHIC (Ref [19])

Strong-strong simulations using the code BBDEMO2C were also done with the appropriate beam parameters, reference [19]. Figure 19 shows a comparison of the measured and simulated bunch spectrum. Without collisions, no π mode was visible but the mode appeared when the
beams were colliding. The measured frequency split between the $\pi$ mode and the $\sigma$ mode was $\nu(\pi) - \nu(\sigma) = 1.3\xi$, in good agreement with the Yokoya parameter 1.31. The coherent mode frequencies calculated by BBDEMO2C were also in good agreement with the observed values.

\[ \pi \] mode \hspace{1cm} \sigma \] mode

\[ \text{Equal tunes} \]

\[ \text{Split tunes} \]

Figure 20: Bunch spectrum calculated with Beambeam3D showing the absence of the $\pi$ mode when the tunes are split by more than $\xi$ (courtesy: J. Qiang)

These coherent modes have not hurt machine performance so far. Usually a small tune split between the beams suffices to remove the $\pi$ mode. This was predicted by A.Hofmann and has also been observed in simulations. Figure 21 shows the bunch spectrum calculated with the code Beambeam3D in two situations. In the first figure, both beams have the same tunes and in the second the tunes are moved apart by 0.02 while $\xi = 0.004$.

VII. LHC Beam-beam simulations

The LHC will be in a new regime where both long-range interactions and strong-strong effects will play a role. Considerable effort has been invested in simulations and analysis to understand the potential limitations. A comprehensive overview can be found in Reference [20]. The areas of effort include:

- Dynamic aperture due to beam-beam interactions and magnet field errors
- Orbits and tunes along the bunch trains
- Proper choice of crossing planes
- Beam-beam compensation with wires
- Impact of ground motion on emittance growth
- Impact of luminosity monitoring on emittance growth
- Excitation of synchro-betatron resonances
- Halo generation
- Possible loss of Landau damping and its restoration
- Alternative paths towards higher luminosity

Reference [20] should be consulted for more details about these studies. Here I’ll focus on a few topics.

The weak-strong code WSDIFF has been used to examine diffusion generated by beam-beam interactions in several situations.

Figure 21: Diffusion rate with nominal beams and during commissioning with different bunch spacings and $\beta^*$ values, taken from Reference [20].
One example, seen in Figure 21, compares the diffusive dynamic aperture during nominal operation with that during commissioning with lower intensity bunches, different bunch spacing and higher $\beta*$ values. At nominal parameters there is a sharp increase in the diffusion rate at $6\sigma$ while with the commissioning beam and 75nsec bunch spacing there is no such sharp increase. This predicts that the limitations due to the beam-beam interactions will not be severe during commissioning with these parameters. Another example shows the use of simulations in choosing parameters. The nominal plan calls for the beams crossing in the vertical plane at IP1 and in the horizontal plane at IP5. This results in a cancellation of the long-range tune shifts and perhaps other benefits. However calculation of the diffusive aperture with WSDIFF, seen in Figure 22, shows larger apertures when the crossing planes are chosen to be the same at both IPs. This topic is still under active investigation.

![Figure 22: Diffusive dynamic aperture with different crossing schemes, Reference [20].](image)

Compensation of the long-range interactions with current carrying wires is also under active investigation. In addition to simulation studies, beam studies have also been done at the SPS [21, 22]. Simulations with WSDIFF shown in Figure 23 found that a wire leads to a similar sharp diffusive aperture as due to the long-range interactions. These results were qualitatively confirmed during the first set of studies in the SPS.

![Figure 23: Diffusion rate due to a wire acting on a single beam, Reference [20].](image)

![Figure 24: Alignment tolerance of the 2nd wire relative to the 1st wire – July 2004 studies [22]. Top: Measured intensity loss vs. the relative vertical displacement between the wires. Bottom: Calculated beam lifetime with BBSIM vs. the relative vertical displacement.](image)
In a later set of studies performed in July 2004, the effect of one wire on the beam was compensated by another wire [22]. The alignment tolerance of the 2nd wire relative to the 1st wire was one of several measurements. In addition, simulations with the code BBSIM were also done prior to the experiments. The top plot in Figure 24 shows the measured beam intensity loss as the relative vertical alignment was changed. The smallest beam loss occurred with $\Delta y = 1$ mm and not at $\Delta y = 0$, possibly due to a mis-calibration. The experiment showed that the compensation was not effective when the 2nd wire was offset by more than 3 mm from the optimum position. The simulation showed that the beam lifetime had the same value (~ 2hrs) when the 2nd wire was offset by more than 2 mm as when the 2nd wire was absent, i.e. there was no compensation with an offset $\geq 2$mm. Thus measurements and simulations were roughly consistent.

There has been a concern that the coherent modes in the LHC excited by the beam-beam interactions may not be Landau damped. The bunch spectrum has been calculated by several simulation codes with different approximations. More recently a 6D parallelized code BEAMX was used to calculate the bunch spectrum for different crossing angles [11]. As seen in Figure 25, this simulation predicts that the $\pi$ mode will be inside the continuum of modes at the full crossing angle of 300 $\mu$rad and probably not a source of concern.

VIII. SUMMARY

Beam-beam phenomena in hadron colliders have been studied for several decades now, yet new phenomena and limitations are observed as existing colliders are pushed to new regimes or new colliders are commissioned. The impact of long-range interactions has been strongly felt during Run II in the Tevatron. We have learnt again that beam-beam performance cannot be characterized by the beam-beam tune shift alone. Several other parameters such as the chromaticity, beam separations around the ring, mismatched emittances as well as the usual suspects: tunes and coupling, determine how the beam-beam interactions limit beam quality. Until recently simulations were mostly used to calculate quantities that are not often measured such as dynamic apertures and tune shifts with amplitude. However the wide availability of parallel computers and fast algorithms now make calculations of emittance growth and lifetimes feasible in many situations. Experience has shown that it is important to include an accurate fully six dimensional model of the linear optics in the simulation model. We still have to make judicious choices about the other beam physics effects (rest-gas scattering, intra-beam scattering, magnetic nonlinearities, impedances, power supply ripple and fluctuations etc.) to include in the simulation model. The simulation results presented here for the Tevatron, RHIC and the SPS are consistent with measurements in some cases. The next step for the case of Tevatron simulations will be to demonstrate a qualitative agreement between simulated bunch-by-bunch lifetimes and emittances and measurements. Beam-beam simulations will have come of age when we can

Figure 25: Bunch spectrum with the strong-strong code BEAMX showing the $\pi$ mode at different crossing angles, Reference [11].
use them routinely to improve the performance of a collider.

ACKNOWLEDGEMENTS


REFERENCES

[8] Tevlat, written by A. Russell and N. Gelfand, FNAL
[22] J.P. Koutchouk et al. to be published
Abstract

I present a brief summary of my remarks made at the Panel Discussion on Beam-Beam Simulations (Session 4A) held at the First CARE-HHH APD Workshop on Beam Dynamics in Future Hadron Colliders and Rapidly Cycling High-Intensity Synchrotrons “HHH 2004” (CERN, 8-11 November, 2004).

Preliminary Questions

It does not seem easy, in general, to benchmark beam-beam simulations against experiments, whether for hadron colliders or for $e^+e^-$ colliders. The beam-beam interaction is almost by necessity designed to be weak, and effects other than beam-beam can obscure the interpretation of the measurements. A faithful simulation, taking into account all relevant effects, is not yet at hand. Nevertheless, weak-strong simulations provide valuable information on the tune footprint.

A computer simulation for beams with many bunches requires large memory and parallel processing. Complete simulations with thousands of bunches are probably unrealistic for the next ~2 – 4 years. At present, strong-strong simulations of the LHC with a few bunches and one IP only are possible at the NERSC computer center at LBNL employing up to 512 parallel processors using 1-2 TB of memory and lasting for ~12 h wall-clock time. Such simulations can track two colliding beams with one bunch per beam with ~ $10^9$ macroparticles per bunch for ~ $10^6$ turns, in linear-lattice approximation. At high intensities self-consistent simulations are likely needed if one wants to answer long-time effects such as emittance growth. This may be especially so for nonzero crossing angle and off-set (parasitic) collisions. Similar comments apply to the necessity of 6D simulations.

Additional Issues

Long-term $\epsilon$ growth.

A realistic strong-strong simulation of the LHC for $10^9$ turns (which is ~25 h LHC time, roughly the expected storage time, and roughly equal to the damping time at top energy) is computationally unrealistic for the foreseeable future. However, strong-strong simulations for $10^6$ turns (~90 s LHC time) and $10^6$ macroparticles per bunch have been done with the code BEAMBEAM3D at NERSC[1]. Such simulations are based on a simplified machine model, namely one IP, one bunch per beam, linear lattice optics, and head-on collisions. Typically 32 computer processors are used, and the simulation takes ~10 h wall-clock time. By employing more processors1 or other computers such as Red Storm at ORNL, one can envision at present simulations for $10^7$ turns (~15 min LHC time) with a few (~3) bunches/beam. Such capability allows beginning to assess emittance growth rate, orbit squeeze, and parasitic collisions in the strong-strong regime. However, it seems important to carefully plan such simulations for maximum profit. We welcome any advice on exactly what simulations would be most useful.

Modeling issues.

As beam-beam simulations become numerically more precise owing to computer power increases, the traditional techniques for such calculations must be re-examined for subtle effects that have heretofore remained weak. This re-examination is especially important when attempting to compute very slow, and probably small, emittance growth in hadron colliders. Four such subtle effects come to mind. It is possible that there are others, or that some answers already exist for some of these:

Binary collisions. In the traditional approach to beam-beam simulations one computes the kick on a macroparticle due to the collective electromagnetic (EM) field of the opposing bunch. However, binary (direct) proton-proton collisions must have an effect at some level of detail. In a faithful simulation, one would represent both colliding bunches by macroparticles of unit charge (ie., one per proton), and one would compute the EM fields form these (or one could carry out a direct-collision simulation). Such a simulation is beyond present-day computer capabilities and, for most practical applications, would be unnecessarily precise. In present-day simulations, each macroparticle represents a large number of protons, typically $10^5 - 10^6$. Furthermore, the EM field produced by these is smoothed in some fashion before the kick is applied to the opposing bunch. Both approximations have the effect of weakening the binary collision effects. It would be valuable to study these two approximations in a methodical fashion in order to quantify their effect on long-term, slow, emittance growth.

Numerical collisionality. Computational approximations lead to purely numerical emittance growth. It is important to quantify this effect, and perhaps minimize it, in order to extract physical results from simulations. An example can be seen in Fig. 1. In another example, carried

\footnote{1The NERSC machine currently has 8000 processors.}
out with a Gaussian strong-strong code for LHC parameters in the impulse approximation, head-on collisions and $10^{12}$ particles per bunch, exhibits the scaling rule $\sigma(n) \propto M^{-p}$ for the transverse rms bunch size $\sigma$ after $n$ turns. The power $p$ was found to be $p \approx 0.7 \sim 0.8$, but the calculations were done only for $M = 10^2$, $10^3$ and $10^4$ and $n = 10^4$ turns [3].

Figure 1: Emittance growth due to the macroparticle charge [2]. Plotted is the scaled emittance, $\epsilon(n)/\epsilon(0)$, at the end of $n = 10^6$ turns, for a strong-strong LHC simulation as a function of macroparticles per bunch $M$. The red curve is the average of the $x$ and $y$ emittances of beam 1, while the green is the corresponding quantity for beam 2. $M$ was the only parameter that was varied from simulation to simulation, from $2 \times 10^5$ to $2 \times 10^6$.

**Space-charge.** It is possible that the space-charge force will have to be included in long-term simulations. The space-charge tune spread $\Delta \nu_{\text{sp.ch.}}$ is $\propto 1/\gamma^2 = 2 \times 10^{-5}$. Even though this is a small number, in the most simplistic estimate it is multiplied by the number of turns, yielding $T/\gamma^2 = 20$, a rather large number (we assume one damping time, $T = 10^9$ turns). Another way to argue is that $\Delta \nu_{\text{sp.ch.}} \propto C/\gamma^2$ where $C$ is the circumference. For $T$ turns, the effective circumference is $TC$, hence the above estimate.

**Beam divergence.** Typical (perhaps all?) simulations for head-on collisions (i.e., zero crossing angle) make the approximation that the particles move collinearly during the beam-beam collision. This is not exact: the particles have a typical angular divergence $\sigma^* = \epsilon/\sigma^*$, which is $3 \times 10^{-5}$ rad for the LHC at collision. This angular divergence leads to a nonzero longitudinal component of the magnetic force $\mathbf{v} \times \mathbf{B}$, which is neglected in conventional beam-beam simulations. This phenomenon is related to, but not equivalent, to a crossing angle.

**ACKNOWLEDGMENTS**

I am grateful to Andreas Adelmann and Ji Qiang for discussions.

**REFERENCES**


RF coupling impedance measurements versus simulations

A. Mostacci∗, L. Palumbo, Università La Sapienza, Roma, Italy
B. Spataro, LNF-INFN, Frascati, Italy,
F. Caspers, CERN, Geneva, Switzerland.

Abstract

Bench measurements nowadays represent an important tool to estimate the coupling impedance of any particle accelerator device. The well-known technique based on the coaxial wire method allows to excite in the device under test a field similar to the one generated by an ultra-relativistic point charge. Nevertheless the measured impedance of the device needs comparisons to numerical simulations and, when available, theoretical results. We discuss the basics of the coaxial wire method and report the formulae widely used to convert measured scattering parameters to longitudinal and transverse impedance data. We discuss typical measurement examples of interest for the LHC. In case of resonant structures, impedance measurements and comparison with simulations become easier. The bead-pull technique may be used in this case.

INTRODUCTION

The interaction between a (relativistic) beam and its surroundings is usually described in terms of longitudinal and transverse coupling impedance [1]. The longitudinal impedance accounts for the energy lost by a point charge $q$ because of the wake field of a leading particle; assuming an infinitely long pipe, for a relativistic beam it is defined as

$$Z_L(\omega) = -\frac{1}{q} \int_{-\infty}^{\infty} E_z(r = 0; \omega) \exp\left(\frac{j \omega}{c} z\right) dz,$$  \hspace{1cm} (1)

where $E_z$ is the longitudinal electric field and $c$ is the speed of light. Instead the transverse kick experienced by a particle because of deflecting fields excited by a leading charge, can be described in terms of the transverse coupling impedance

$$Z_T(\omega) = \frac{j}{q^2} \int_{-\infty}^{\infty} F_T(r_1, r_2; \omega) \exp\left(\frac{j \omega}{c} z\right) dz,$$  \hspace{1cm} (2)

where $F_T$ is the transverse Lorentz force and $r_1$ ($r_2$) is the leading (trailing) particle position. Longitudinal impedance is, therefore, measured in $\Omega$ while the transverse one in $\Omega/m$.

Often beam dynamics in circular machines is studied in the frequency domain due to the intrinsic periodicity and then coupling impedances are used to estimate the effect of any accelerator device on the beam stability. Coupling impedance is then an important design parameter for any element to be installed in the machine. In the LHC case, the impedance of most elements is carefully optimized through numerical simulations, benchmark measurements and confirmed by beam measurements, when available. Coupling impedance data of each LHC device could be stored in a database to eventually compute bunch instability thresholds for any possible machine settings [2].

In the design stage, coupling impedance can be estimated either by numerical simulations or by bench measurements. Nowadays numerical simulations greatly exploit commercial general purpose codes solving Maxwell equations in the Device Under Test (DUT), i.e. computing the electromagnetic field in any point of the structure (3D simulations). Time domain codes exist (MAFIA [3] and GdfidL [4]) directly providing wake-potentials which can be Fourier transformed to get coupling impedances. Other general purpose codes can be, as well, used to estimate coupling impedance of any DUT by simply simulating the bench measurement discussed below; for example that is the case of HFSS [5] (frequency domain) or MWstudio [3] (time domain). Less general codes, studying EM fields in symmetric (2D) structures (e.g. SUPERFISH [6] or OSCAR2D [7]) are used as well. There exist also “dedicated” codes which compute only the longitudinal and transverse coupling impedances of 2D structures in a very reliable way; the most famous is ABCI [8]. Numerical codes computing coupling impedances in specific geometries have also been developed, but they will not be discussed here.

The aim of this paper is to briefly review the standard way of bench measuring coupling impedances and to discuss comparisons with numerical simulation, through some meaningful examples. The coaxial wire method is the most well known bench method and it is firstly discussed for longitudinal and transverse impedance measurement. The wire set-up can also be used to study the properties of the structure when excited by a beam passing through; trapped modes or beam transfer impedance can, for example, be measured in this way. Impedances in resonant structures (e.g. accelerating or deflecting cavities) deserve a different treatment and they are measured with bead-pull techniques, discussed in a later section. Throughout the paper, we will try to discuss the agreement between measurement and simulations in typical cases, pointing out some open questions as well.

THE COAXIAL WIRE METHOD

Motivation and validation

The field of a relativistic point charge $q$ in the free space (or in a perfectly conducting beam pipe) is a Transverse Electric Magnetic (TEM) wave; namely it has only compo-
nents transverse to the propagation direction (z-axis):

\[ E_z(r, \omega) = Z_0 q \frac{Z_0 q}{2\pi r} \exp \left( -j \frac{\omega}{c} - \frac{1}{c} \right), \quad (3) \]

being \( E_z \) (\( H_z \)) the electric (magnetic) radial (azimuthal) field component. The amplitude scales inversely with the distance \( r \) from the propagation axis and the propagation constant is \( \omega/c \). The fundamental mode of a coaxial wave guide is a TEM wave as well, with the same amplitude dependence on \( 1/r \) and the same propagation constant. For example, let us consider a perfectly conducting cylindrical wave guide with circular cross section of radius \( b \) with a conductor of radius \( a \) on the axis of the guide. The resulting coaxial structure allows the propagation, at any frequencies, of a TEM mode whose field is given by

\[ E_z(r, \omega) = Z_0 A \frac{A}{r} \exp \left( -j \frac{\omega}{c} z \right), \quad (4) \]

where \( A \) is a constant depending on the power actually flowing in the guide. Therefore the excitation due to a relativistic beam in a given DUT can be “simulated” by exciting a TEM field by means of a conductor placed along the axis of the structure. The impedance source on the DUT will scatter some field, i.e. exciting some higher order modes; such modes must not propagate otherwise the propagating field will not be anymore similar to the TEM beam field. In principle, then, simulating the beam field with the TEM mode of a coaxial waveguide is possible only at frequencies below the first higher mode cut-off, namely below the TM01 cut-off frequency. One can also demonstrate that the modes of the coaxial waveguide converges for vanishing wire radius to the analogous mode of the cylindrical waveguide, at least at the beam pipe boundary, where the impedance source is usually located [9].

To compare the excitation of a given DUT by a coaxial wire and with the beam itself, we are going to discuss some measurements done in the framework of the investigations of the shielding properties of coated ceramic vacuum chambers [10]. The 500 MeV CERN EP A electron beam was sent through two identical ceramic vacuum chamber sections; the first one was internally coated with a layer of 1.5 \( \mu \)m depth (DC resistance of 1 \( \Omega \)). Magnetic field probes were placed to measure the beam field just outside the two ceramic chambers (the coated and the reference one). In 1999 experiment, shielding properties of the resistive coating (thinner than the skin depth) were demonstrated, confirming previous indirect measurements and simulations [11]. In the 2000 experiment, among other results, it was proved that the screening properties of the coating can be spoiled by the addition of a second conducting layer placed outside the field probes and electrically connected to the metallic vacuum chamber sections. In this case, in fact, the magnetic field probe was measuring clearly the field of the 1 ns (r.m.s.) bunched beam (see Fig. 1). The same chamber in the same configuration (i.e. with this additional external conductor) was then measured in the bench set-up: a 0.8 mm diameter wire was stretched on the axis of the structure. One end of the wire was connected to a 50 \( \Omega \) load while the other end was connected to one port of a Vector Network Analyzer (VNA); matching resistors were used. The other port of the VNA was connected to the field probe. The network analyzer was set to send through the wire a synthetic pulse (time domain option) with 300 MHz bandwidth and measured the transmission between the ports, i.e. the signal through the probe. This particular kind of set-up is not very often used, but it is very similar to the “time domain” measurement originally proposed by Sands and Rees in the 70s’ [12]; nowadays time domain measurements are often performed with synthetic pulse techniques in many microwaves applications. The measurement with the beam and with the wire should give virtually the same result, apart from a scaling factor due to the difference of the power carried by the beam and by the VNA signal. The results are shown in Fig. 1 where the beam and the bench data have been normalized and time shifted so that the traces coincides in their minimum point. The external shield, having a DC resistance much smaller than the coating, carries the image currents, the field penetrates the ceramic and the field probe can measure a clear signal. This is only one of the configurations measured both with the beam and in the bench set-up; the agreement with other measurements is similar to the one of Fig. 1. The results of that comparison confirm the validity of the coaxial wire approach to simulate the beam field effect on a given DUT.

![Figure 1: Signal from the field probe after normalization and time shifting in the EPA experiment on coated chamber shielding properties. The field probe is inserted between the coated ceramic and an external conductor connected to the beam pipe.](image)

The coaxial wire bench measurement exploits the analogy between a relativistic point charge field and the fundamental mode of a coaxial waveguide. Therefore a metal wire is stretched in the DUT which is then transformed in a coaxial transmission line. Thus the longitudinal coupling impedance can be inferred from the properties of such a transmission line, provided that only the (fundamental) TEM mode is propagating. To bench measure longitudinal
impedance, a wire along the axis of the structure is needed, while transverse impedance can be measured from a single displaced wire or two wires placed symmetrically with respect to the axis. A detailed review of the method, including some practical measurement tips, can be found in Ref. [13].

**Longitudinal Coupling Impedance**

The wire stretched in the DUT of length $L$ can be modeled, as mentioned before, as a TEM coaxial line of length $L$. In general, such a line is considered to have distributed parameters but in case of $L$ much smaller than the wavelength $\lambda$ the lumped elements approximation is applicable. The DUT beam coupling impedance is then modeled as a series impedance of an ideal reference line (REF). Therefore coupling impedance can be obtained from the REF and DUT characteristic impedances and propagation constants of the lines (see for example Ref. [14]). It is well known that any transmission line can be characterized by measuring its scattering $S$-parameters, for example with VNA. In principle both reflection (i.e. $S_{11}$) and transmission measurement (i.e. $S_{21}$) are possible, but usually transmission measurement is preferred for practical reasons. In the framework of this transmission line model, the DUT coupling impedance can be exactly computed from measured $S$-parameters but the procedure is cumbersome and not practically convenient. Therefore a number of approximated formulae are derived in literature and we will report the most used ones, highlighting their approximations. All the following formulae do not consider the effect of the mismatch at the beginning and at the end of the perturbed transmission line. Therefore matching networks (resistive networks or cones) are normally used in the actual bench set-ups. Cones are mechanically difficult and act as a frequency dependent distributed transformer which doesn’t work at low frequency; on the contrary resistive networks are affected by parasitic inductances and capacitances affecting their performance at high frequency (depending on the components actually used). Approximated formulas and the “exact” transmission line solution are numerically compared in Ref. [15].

Being $Z_c$ the characteristic impedance of the wire inside the DUT, the beam coupling impedance $Z_{||}$ can be estimated with the “improved log-formula” [14]:

$$Z_{\text{LOG}} = -Z_c \ln \left( \frac{S_{21}^{\text{DUT}}}{S_{21}^{\text{REF}}} \right) \left[ 1 + \frac{\ln \left( \frac{S_{21}^{\text{DUT}}}{S_{21}^{\text{REF}}} \right)}{\ln \left( \frac{S_{11}^{\text{DUT}}}{S_{11}^{\text{REF}}} \right)} \right]. \tag{5}$$

Expressing the $S_{21}^{\text{REF}}$ in terms of the DUT electrical length $L$ one can get another equation analogous to Eq. (5) [16]:

$$Z_{\text{LOG}} = -Z_c \ln \left( \frac{S_{21}^{\text{DUT}}}{S_{21}^{\text{REF}}} \right) \left[ 1 + \frac{j c}{2 \omega L_f} \ln \left( \frac{S_{21}^{\text{DUT}}}{S_{21}^{\text{REF}}} \right) \right], \tag{6}$$

which can be useful in practice. The improved impedance expression requires the knowledge of the electrical length of the DUT and its accuracy decreases for shorter devices [17]. Reference [15] suggests the use of improved log-formula for DUT longer than the wavelength $\lambda$.

For small ratios $Z_{||}/Z_c$ the so called “standard log-formula” has been proposed for the distributed impedances [18]:

$$Z_{\text{log}} = -2Z_c \ln \left( \frac{S_{21}^{\text{DUT}}}{S_{21}^{\text{REF}}} \right). \tag{7}$$

The log-formula Eq. (7) is generally applicable including lumped components, provided that no strong resonance is present and the perturbation treatment is justified.

For lumped elements, i.e. when the DUT electric length is much smaller than the wavelength, the previous expressions converge to the so called “lumped element formula” [19]:

$$Z_{\text{HP}} = -2Z_c \frac{S_{21}^{\text{DUT}} - S_{11}^{\text{DUT}}}{S_{21}^{\text{REF}}}. \tag{8}$$

The lumped impedance formula is applicable to single resonances and has the advantage that the scattering coefficient ratio is directly converted into an impedance by the network analyzer [20].

The quantities $Z_{\text{LOG}}$, $Z_{\text{log}}$, $Z_{\text{HP}}$ are estimations of the beam coupling impedance $Z_{||}$; the smaller the ratio $Z_{||}/Z_c$, the more accurate are the approximated formulae. As an example, a wire in a perfectly conducting cylindrical beam pipe with circular cross section has a characteristic impedance equal to

$$Z_c = \frac{Z_0}{2\pi} \ln \left( \frac{b}{a} \right) \tag{9}$$

where $a$ is the wire radius, $b$ is the (inner) pipe radius and $Z_0$ is the vacuum impedance. Therefore a smaller wire has an higher $Z_c$, resulting in a more accurate measurement of the coupling impedance. A detailed discussion of the systematic error done in estimating the beam coupling impedance $Z_{||}$ with $Z_{\text{LOG}}$, $Z_{\text{log}}$ or $Z_{\text{HP}}$ is reported in Ref. [17].

The difference between the improved log formula Eqs. (5, 6) and the standard one Eq. (7) can be shown in measurements performed on the 7 cells module of the MKE kicker [23]. The coupling impedance is much bigger than the characteristic impedance of the wire in the DUT ($\approx 300\Omega$) and therefore the improved log formula must be used:

$$Z_{\text{LOG}} = Z_{\text{log}} \left[ 1 + \frac{j c}{2 \omega L_f} \ln \left( \frac{S_{21}^{\text{DUT}}}{S_{21}^{\text{REF}}} \right) \right]. \tag{10}$$

Equation (10) differs from Eq. (6) because the length of the ferrite ($L_f = 1.66m$) is used instead of the length of the whole kicker tank ($L = 2.31m$), as discussed in Ref. [23]. Figure 2 shows the wire measurement results interpreted with the improved formula Eq. (10) (green line) and the standard one Eq. (7) (blue line). The comparison with theory (black line) shows that, at least for the
real part of the impedance, the improved log formula gives a result closer to theoretical expectations for frequencies higher than few hundreds of MHz. At lower frequencies, i.e. where the DUT length is comparable to the wavelength and the impedance is much closer to the characteristic impedance of the wire in the DUT, the standard log formula is a better estimation of the coupling impedance.

As an example of a comparison between numerical simulation and bench measurement of longitudinal coupling impedance, we briefly discuss the case of the LHC injection kicker. In particular, to study the impedance of the LHC injection kicker a 1 meter long model has been built and bench measured [21]; the measured impedance data have been also compared to HFSS simulations [22].

The model consists of a ceramic test chamber (shown in Fig. 3) with 30 printed conducting strips inside; the strips have different widths and are done using the same technology of the final LHC kicker. Good RF-contacts are assured on either side (“A” and “C” in Fig. 3). The lower part (rectangular steel profile underneath the ceramic tube) serves both as mechanical support and as electrical simulation for the “cold conductor” of the kicker. The “printed” coupling capacitor is visible near the right connector (position “C”) and has pyramid like “teeth” toward position “B”. These “teeth” are intended as an RF match (gradual transition for the image current) toward higher frequencies. A coaxial Cu-Be wire is stretched on the axis of the structure and it is soldered through matching resistors to the two RF connectors in each side of the chamber. Such a bench measurement set-up has been simulated with HFSS, including the current bypass conductor, but with 12 identical strip lines of the same width to avoid complexity [22].

We refer for the details of the measurement to Ref. [21] and for the simulation to Ref. [22], where more extensive explanations can be found. The coupling impedance exhibits some peaks at particular frequencies being negligible in all the rest of the measured frequency range. To compare measurement and simulations, we have reported in Table 1 the values of the main peaks of the longitudinal coupling impedance measured with the wire bench set-up and estimated with the HFSS simulation of the wire measurement itself. The agreement between measurement and simulations shown in Table 1 is satisfactory considering that only 12 over 30 conducting strips are included in the simulations, which essentially explains the difference.

According to Ref. [22], the origin of the 17 MHz peak is the resonance of the capacitor at one end and the inductance created by the strips and the outer support. To damp this low frequency resonance a lossy ferrite ring was inserted; measurements and simulations agree in estimating a reduction of the amplitude of the impedance. The 31 MHz peak is mainly due to a transverse resonance and it is very much affected by an imperfectly tightened wire or by a small offset of the wire itself. In the simulations [22] this resonance is practically canceled with a 2 mm offset of the wire; measurements give similar results [21]. The origin of the 442 MHz and of the 846 MHz peaks is by the coaxial waveguide resonance at the copper tape (point “C” in Fig. 3).

In conclusion, longitudinal coupling impedance bench measurements are reasonably well understood and the technique is well established. With modern simulation codes, one can derive directly the coupling impedance or simulate the bench set-up with wire, virtually for any structure. Evaluation of coupling impedance from measured or simulated wire method results require the same cautions; but simulations and RF measurements usually agree well. Moreover comparison with numerical results are very useful to drive and to interpret the measurements. One should pay attention that simulation may require a simplified DUT model which will only reproduce the main DUT electro-

Figure 2: Real part of the longitudinal coupling impedance for the 7 cells MKE kicker module [23]. The measured data are interpreted via the improved log formula (green line) or the standard log formula (blue line) and compared to theoretical expectations (dashed line).

Figure 3: Model for the LHC injection kicker [21].

<table>
<thead>
<tr>
<th>Simulations [22]</th>
<th>Measurements [21]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Freq. (MHz)</td>
<td>(Z_{\parallel} (\Omega))</td>
</tr>
<tr>
<td>(Z_{\parallel} (\Omega))</td>
<td>Freq. (MHz)</td>
</tr>
<tr>
<td>14</td>
<td>500</td>
</tr>
<tr>
<td>28</td>
<td>44</td>
</tr>
<tr>
<td>410</td>
<td>65</td>
</tr>
<tr>
<td>810</td>
<td>40</td>
</tr>
</tbody>
</table>
magnetic features.

Transverse Coupling Impedance

The transverse impedance is proportional at a given frequency to the change in longitudinal impedance due to the lateral displacement of the beam in the plane under consideration (vertical or horizontal). Therefore the transverse impedance is proportional to the transverse gradient of the longitudinal impedance (Panowsky-Wentzel theorem [1]).

Based on this theorem, the most common method to bench measure transverse impedance uses two parallel wires stretched along the DUT [13]. Opposite currents are sent through the wires (odd mode excitation); instead of the wires, a loop can be used to increase signal to noise ratio [25]. The bench transverse impedance $Z_{\perp,\text{bench}}$ is given by [24]

$$Z_{\perp,\text{bench}} = \frac{Z_{\parallel,\text{bench}}C}{\omega \Delta^2},$$

where $\Delta$ is the wire spacing (usually about 10% of the DUT radius). $Z_{\parallel,\text{bench}}$ is the longitudinal coupling impedance measured from the $S$-parameters as discussed above, e.g. using the improved log-formula

$$Z_{\perp,\text{LOG}} = \frac{Z_c}{\omega \Delta^2} \ln \left( \frac{S_{21}^{DUT}}{S_{21}^{REF}} \right) \left[ 1 + \ln \left( \frac{S_{21}^{DUT}}{S_{21}^{REF}} \right) \right],$$

where now $Z_c$ is the characteristic impedance of the odd mode of a two wires transmission line. Concerning LHC (and other future machines as well), low frequency transverse impedance is interesting and therefore the lumped element Eq. (8) must be used in Eq. (11). A practical example of low frequency transverse impedance is reported in Ref. [25] for a simple case; results are compared to theoretical expectations to define a reliable measurement procedure.

In the two wires bench set-up only dipole field components are excited because of the symmetry of the wires/coil; therefore there is no electric field component on the axis. In numerical simulations, this is analogous to putting a metallic image plane between the wires. Nevertheless some accelerator devices may exhibit a strong asymmetry in the image current distribution due to azimuthal variation of conductivity (e.g, ferrite in kickers) or to cross section shape. Two wire techniques can be used with some caution in this cases because the field in the structure is not TEM-like; in order to get a more complete view of the transverse kick on the beam, it may be useful to characterize the device with a single wire [26].

The transverse impedance itself can be measured with a single wire displaced in various positions, that is measuring the longitudinal coupling impedance as a function of the displacement $x_0$ of a single wire. From the Panowsky-Wentzel theorem we get

$$Z_{\perp,\text{bench}} \approx \frac{1}{\omega} Z_{\parallel,\text{bench}}(x_0) - Z_{\parallel,\text{bench}}(x_0 = 0),$$

provided that $x_0$ is small with respect to the typical variation length of the bench measured coupling impedance $Z_{\parallel,\text{bench}}$.

From the practical point of view, transverse impedance measurement techniques are more delicate and require particular attention for asymmetric devices (e.g. traveling wave kickers like). Novel techniques optimized for particular DUTs, are also being proposed, e.g. SNS kicker measurements reported in Ref. [27]. Numerical simulations are necessary to control and validate the measurement procedure. One should pay attention that DUT models feasible for simulations do not introduce non physical symmetries or approximations; in principle, dealing with transverse problem may require more complex simulations than the longitudinal case.

Other applications: trapped mode finding

A coaxial wire set-up can also be used to study the behavior of a given DUT when traversed by a relativistic beam. For instance, it can be useful to measure the beam transfer impedance or to check if trapped mode are excited in a given structure. The beam transfer impedance is the ratio between a voltage signal induced because of the beam structure interaction and the beam current. Examples of beam transfer impedance for a bunch length monitor is reported in Ref.s [29, 30], while we are going to discuss briefly how a coaxial wire set-up has been used to study the possibility of trapped modes in the LHC recombination chamber.

In the LHC interaction regions, the separated beam pipes join in a common region to allow collisions. Thus the beam “sees” an enlargement of the beam pipe that, if not properly tapered, can lead to trapped modes dangerous for the beam stability. Such modes exist in the LHC chambers, but luckily enough they are too weak to affect the beam stability. To prove it, numerical simulations and bench measurements were done as reported in Ref. [28] and related references where the interested reader can find additional details. In the following we report part of that work, in particular focusing on comparisons between measurements and simulations.

A simplified (rectangular and scaled) model of the recombination chamber was built to be bench measured with a coaxial wire set-up. In the original design the transition is tapered; in our model the taper was first removed to prove the existence of the trapped mode and then added again to show the reduction of the mode amplitude.

To demonstrate the feasibility of the measurement, HFSS simulations were done and the results are shown in Fig. 4. Picture a) of Fig. 4 shows the electric field at 2.737 GHz: the trapped mode at the wedge is excited and there is no transmission along the wire: the signal is reflected and one can see a stationary wave pattern of the field. Picture b) of Fig. 4 shows that at other frequencies (e.g. 2.5 GHz) the field is confined around the wire but can propagate from one end of the wire to the other. There-
fore, the HFSS simulations show that the trapped mode is present and that it can be measured by looking for a strong notch in the transmission coefficient between the two ends of the wire.

Figure 4: Amplitude of the electric field (maximum value) for the bench set-up of the LHC recombination chamber at 2.737 GHz (left plot) and 2.5 GHz (right plot) by HFSS simulations.

The same geometry, i.e. without tapering, have also been simulated with MAFIA, i.e. computing directly wake potentials. The result for the real part of the coupling impedance \( Z_{RE} \) is shown in Fig. 5. The strong peak at 2.800 GHz is due to the trapped mode.

![Figure 5: Real part of the longitudinal coupling impedance for the bench set-up of the LHC recombination.](image)

Measurements of the transmission between the ends of the wire are shown in Fig. 6 in terms of the transmission scattering parameter \( S_{12} \). As expected, there is a strong notch at 2.753 GHz where the trapped mode is excited and the electromagnetic power remains confined in the structure.

Additional measurements and MAFIA simulations have been done including a taper resembling the actual geometry of the recombination chamber. The beneficial effect of reducing the trapped mode has been shown both from measurement and from simulations [28].

Thus also in this case, measurements and simulations give close results. It is worth noticing that while HFSS gives actually simulations of the bench measurement, MAFIA simulations compute directly the wakefields, that is the structure is “excited” directly by a particle beam; the agreement of the results, therefore, assesses the coupling impedance measurement and computation techniques.

### RESONANT STRUCTURES

An important class of accelerator devices are cavities which are now used both for accelerating and deflecting the particle beam. Each cavity is characterized by its resonant frequency \( f_0 \), the quality factor of the resonance \( Q \) and its shunt impedance \( R \). One can think of measuring all these quantities with the coaxial wire set-up, i.e. measuring strong notches in the transmission coefficient between the ends of the wire. But the wire perturbs longitudinal cavity modes, e.g. lowers the \( Q \) and detunes the frequency. Therefore the coaxial wire set-up is not usually recommended for cavity measurements and it is advisable only for special cases, mainly transverse modes [13].

The most used technique to characterize cavities is the “bead pull” measurement [31]. The field in the cavity can be sampled by introducing a perturbing object and measuring the change in resonant frequency: where the field is maximum (minimum) the resonance frequency will be more (less) perturbed. It is a perturbation method, therefore the perturbing object must be so small that the field does not vary significantly over its largest linear dimension. Shaped beads are used to enhance perturbation and give directional selectivity among different field components.

Quantitatively, the change of the resonant frequency is related to the perturbed cavity field by the Slater theorem. For the typical case of longitudinal electric field on the axis of accelerating cavities, the variation of the resonance frequency \( \Delta f \) from the unperturbed one is [32]

\[
\frac{\Delta f}{f_0} = -\Delta V \epsilon_0 k_E \frac{E_z^2}{4U}
\]

for a conducting bead of volume \( \Delta V \); \( E_z \) is the field at the perturber position and \( U \) is the electromagnetic energy stored in the cavity. The form factor \( k_E \) of the the perturbing object can be exactly calculated for ellipsoids or can be calibrated in a known field (e.g. TM\text{_{001}} of a pillbox cavity).

The frequency variation can be measured by the variation of the phase \( \phi \) at the unperturbed resonant frequency,
according to [33]

$$\Delta f/f_0 = \frac{\tan(\phi(f_0))}{2Q_L} \simeq \frac{\phi(f_0)}{2Q_L}$$

where $Q_L$ is the (loaded) quality factor of the resonance. Even if a very precise initial tuning is needed, this method allows easily measuring the field of many points (as many as the points of the instrument trace). The field shape can also be directly visualized on the instrument screen, greatly facilitating the structure tuning procedure.

As an example we can consider simulations and measurements done on an 11.424 GHz standing wave multi-cell cavity, designed for the SPARC project [34]. The cavity is supposed to work in the $\pi$-mode (all the cells are filled with field) exhibiting a maximum field equal in every cell (field flatness). A 9 cells prototype has been designed and built and all the details are reported in Ref. [34] while a picture is given in Fig. 7.

![Figure 7: Nine cells copper prototype of the SPARC 11 GHz cavity.](image)

Figure 7: Nine cells copper prototype of the SPARC 11 GHz cavity.

Figure 8 compares measured data against numerical simulation results for the electric field on axis. The field has a maximum/minimum in the center of every cell and the tuners have been set to have the required field flatness. The main measurement artifact was the non-negligible effect of the glue used to fix the bead on a plastic wire to be moved by the stepping motor; therefore the glue effect was measured and calibrated away resulting in the data reported in Fig. 8. Numerical codes gives very close results among each other and they all agree well with measurements.

An important cavity design parameter is the $R/Q$ which can be obtained from electric field data using

$$\frac{R/Q}{L_c} = \frac{1}{U \omega L_c} \left| \int_0^{L_c} E_z(z) \exp \left( j \frac{\omega z}{c} \right) dz \right|^2,$$  \hspace{1cm} (12)

where $L_c$ is the length of the structure. From measurement and simulations we get the results given in Table 2 which again shows the good agreement among measurements and simulations. The number in round brackets is the uncertainty of the measurement according to the procedure given in Ref. [35] (Type A evaluation).

Table 2: SPARC 11 GHz cavity: $R/Q$ per unit length.

<table>
<thead>
<tr>
<th></th>
<th>HFSS</th>
<th>SUPERFISH</th>
<th>MAFIA</th>
<th>Meas.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$\Omega/m$</td>
<td>9138</td>
<td>9232</td>
<td>9392</td>
<td>9440(87)</td>
</tr>
</tbody>
</table>

pull measurements are often used to check if the DUT fits the design specifications and they are still required for tuning the multiple cell cavities. $R/Q$ measurements agree always very well with simulations (within the 3%) .

**CONCLUSIONS AND OUTLOOK**

The most common bench methods for measuring coupling impedance have been discussed: wire method (longitudinal and transverse impedance) and bead pull method (resonant structures). There are several codes usable to estimate coupling impedances (MAFIA, HFSS, GdfidL, MWstudio, ABCI, OSCAR2D, SUPERFISH). Impedance can be computed by numerical RF simulators or calculating directly the wakefields or by simulating the wire measurement. Numerical simulations often deal with simplified models: i.e. one must have some insight into the problem. With reasonable and accurate modeling, a longitudinal impedance simulation usually reproduces experimental results (and vice-versa). Transverse impedance requires a much deeper control of both simulations and bench measurements, particularly for some special devices. Impedances of resonant structures are well modeled by numerical simulation with high degree of reliability.

Transverse impedance measurements are surely more challenging than longitudinal ones for modern accelerators devices. Single moving wire techniques should surely be investigated (Panowsky-Wentzel theorem). The measurement techniques strongly depends on the device to be tested; in collimators, for example, one could keep the wire
fixed and move the collimator plates. Nowadays mechanical drawings of complex accelerator components are available in electronic CAD formats and it would be highly desirable to use them as input files of some electromagnetic field simulator. In most cases only simplified schematic models of such complex geometries can be efficiently simulated (e.g., kickers). Moreover electromagnetic models of material used in some devices (e.g., special ferrite in some kickers) are still missing in certain state-of-art commercial numerical electromagnetic field simulators.

ACKNOWLEDGMENTS

The authors are grateful to F. Ruggiero and the CARE-HHH workshop organization team for the invitation to the workshop. H. Tsutsui (Sumitomo) and D. Alesini (INFN) helped a lot with fruitful discussions; T. Linnecar carefully read the manuscript giving helpful comments. The EPA experiment was organized and carried out also by L. Vos, D. Brandt and L. Rinolfi (CERN). V. Lollo, A. Bacci, V. Chimenti (INFN) are involved in the design and the prototyping of the 11 GHz cavity for SPARC. Students of the “Accelerators and detectors” laboratory at Energetics Department of Università “La Sapienza” substantially helped in setting up the bead-pull measurement for SPARC cavities.

REFERENCES

[34] A. Bacci, M. Migliorati, L. Palumbo, B. Spataro, INFN LNF 03/008(R) (2003).
Feasibility of Directly Computing Wakes of Lossy Collimators

W. Bruns, WBFB, Berlin, Germany ∗

RAM-REQUIREMENTS FOR DISCRETISING THE WHOLE GEOMETRY

When one wants to directly compute the wakepotential of a graphite-scrap, one needs a timestep which is smaller than the decaytime within the graphite. The material parameters of graphite, \( \varepsilon_r = 15, \kappa = 0.07 \times 10^6/\Omega/m \), give a decaytime of \( \tau = \varepsilon_0 \varepsilon_r/\kappa = 2 \times 10^{-15}s \). For having a small dispersion error, one wants to have a grid-spacing which is about \( \Delta z \approx C \Delta t = 3 \times 10^8 \times 2 \times 10^{-15}m = 0.6 \times 10^{-6}m \). We therefore want a grid-spacing in the order of \( \Delta z \approx 10^{-6}m \).

The dimensions of a graphite scraper are about 10 cm times 10 cm in the x-y-plane, and the length is in the order of meters.

![Figure 1: The dimensions of the scraper are about 10 cm times 10 cm in the transverse plane, and the length is in the order of meters.](image)

When we directly want to discretise the scraper with a grid-spacing of \( \Delta z = 10^{-6}m \), we need
\[
\frac{V}{(\Delta z)^3} = \frac{0.1m \times 0.1m \times 1m}{(10^{-6}m)^3} = 10^{16}
\]
gridcells. The memory requirement for a gridcell is about 100 bytes. We would need \( 10^{16} \times 10^2 \) Bytes = 10^9 GBytes. Absolutely out of reach. This is probably more than the total amount of RAM of all computers of the world combined.

RAM: SHORT-RANGE WAKEPOTENTIAL

For computing short range wakepotentials up to a length of \( s \), we only need to compute the electromagnetic field in a volume which has a length of \( s \). This volume must move through the geometry together with the exciting charge. With a grid-spacing of \( \Delta z = 10^{-6}m \), we need for computing short range wakepotentials up to \( s = 10^{-3}m \) a total number of gridcells of
\[
\frac{V}{(\Delta z)^3} = \frac{0.1m \times 0.1m \times 10^{-3}m}{(10^{-6}m)^3} = 10^{13}
\]
gridcells. Still much too much.

RAM: ANISOTROPIC MESH

When we have a very short bunch, we need a very small gridspacing in the direction of flight. Within the graphite, we also need a very small gridspacing in the direction of exponential decay. At least near the graphite-vacuum interface. When we use a gridspacing of \( \Delta z = \Delta x = 10^{-6}m \), and a gridspacing of \( \Delta y = 10^{-3}m \), for computing in a moving volume of length \( s = 10^{-3}m \), we need
\[
\frac{V}{\Delta x \Delta y \Delta z} = \frac{0.1m \times 10^{-3}m \times 0.1m}{(10^{-6}m)^3 \times 10^{-3}m} = 10^{10}
\]
gridcells. Now this is within reach. Computing with that number of gridcells will require 500 GBytes of RAM. When the geometry has two planes of symmetry, only a quarter is used. 130 GBytes. Using tricks, one might be able to need only 30 GBytes.

![Figure 2: Because the fields will not vary much as a function of the horizontal coordinate, one can use a grid which has a much larger gridspacing in one direction, than in the other. Fine meshing in the vertical direction is also only required near the vacuum-graphite interface.](image)

* bruns@gdfisl.de
CPU-Time

For computing the wakepotential, the charge and the moving mesh must travel through the whole length of the scraper. As the charge and the mesh cannot travel more than one gridspacing per timestep, we need at least \( \frac{L}{\Delta z} = \frac{1}{10^{-6}} = 10^6 \) timesteps.

The field computation needs about 100 floating point operations per gridcell. The computation over \( 10^6 \) timesteps in a grid of \( 10^{10} \) gridcells then needs \( 10^{18} \) floating point operations. When we have a system capable of \( 1 \text{GFlop} = 10^9 \) Floating point operations per second, the computation would take \( 10^9 \) seconds. This is 280,000 hours. Or 12000 days.

On a cluster of 20 modern PCs, we will be able to achieve 5 to 10 GFlops. On such a system, the computation then will take in the order of 1200 days. This is much too long. We have not yet taken into account, that only a quarter of the structure must be modeled. This would then push the required time down to 300 days. Still much too much.

Inhomogeneous grid in both x- and y-direction is required. Figure 2 sketches a possible grid, where a fine mesh is used only near the vacuum-graphite interface. The total number of gridcells in such a graded mesh will be about \( 1/10 \) of the number without graded mesh in the vertical direction. Therefore the computation time then also will go down by a factor of ten, leaving us with a required time of about 30 days.

**CONCLUSION**

It is possible to directly compute wakes in lossy structure, provided only extremely short range wakepotentials are of interest. One needs a moving mesh on a parallel system. The limiting factor will not be the required memory, but the required computation time. Work is begun to implement such a moving mesh on a cluster of workstations.

Figure 3: Electric field of a gaussian linecharge in a moving mesh a distant times.
Cold to warm transition impedances in the SPS machine.

B. Spataro, D. Alesini, LNF-INFN, Frascati, Italy.
M. Migliorati, A. Mostacci, L. Palumbo, University of Rome ”La Sapienza”, and LNF-INFN, Italy.
F. Ruggiero, Cern, Geneva, Switzerland.

Abstract

In many papers, the interaction between a relativistic particle beam and a vacuum chamber with holes is usually described in terms of coupling impedance and loss factor. The interference among the holes is the main source of wake fields and losses. This note focuses on the impedances evaluation of SPS cold to warm transitions. A comparison between analytical model and numerical results is presented.

INTRODUCTION

In order to set up the SPS machine as final injector for the LHC, the substitution of some components was required. Measurements made with the LHC beam showed that a reduction of the machine impedance was mandatory. In particular in the cold to warm region it has been measured an unacceptable heating (from 0.5 to 4 W/m power loss). For this reason, a detailed study of the coupling impedance budget for both the cold and warm sections was carried out.

Because of the absence of an axial symmetry of the above components, a numerical approach to the problem based on the use of three-dimensional computer codes has been adopted.

To optimize the accuracy of the results and to minimize the CPU time requirements, the cold and warm transitions have been treated separately obtaining independent impedance estimations, even though the two sections represent a unique structure from the impedance point view.

In this work the calculations have been made by supposing copper OFHC structures with a conductivity of $5.91 \times 10^7 \, (\Omega \cdot m)^{-1}$ at room temperature.

The final shape of the cold transition vacuum chamber has a circular cross section and pumping slots for reducing strongly the heating. The warm transition studies gave no specific warning, thus no vacuum chamber modification was needed.

The global cold to warm transition cross section is sketched in Fig. 1.

The geometries of the cold and warm transitions used in MAFIA [1] simulations are reported in Figs. 2 and 3, respectively. Because of the symmetry we have simulated only half structure.

UNIFORMLY SPACED HOLES

The coupling impedance of a circular coaxial beam pipe with $N$ pumping holes has been studied extensively by means of the modified Bethe theory; the most important approximations and the main results are reported in [2, 3]. Being $b$ ($d$) the inner (outer) radius of the coaxial beam pipe and $\alpha_e$ ($\alpha_m$) the electric (magnetic) polarizability of the equally spaced pumping holes, the coupling impedance

Figure 1: Sketch of the cold to warm transition. The left and the right side represent the cold and the warm section respectively.

Figure 2: Shape of the cold transition used for Mafia simulations.

Figure 3: Shape of the warm transition used for Mafia simulations.
reads

\[
Z_{RE}(\omega) = Z_0 \frac{\omega^2 (\alpha_m^2 + \alpha_e^2)}{16\pi b^4 \ln (d/b)c^2} \left[ N + \frac{(\alpha_m + \alpha_e)^2}{\alpha_m^2 + \alpha_e^2} \sum_{h=1}^{N-1} \sum_{k=1}^{h-1} e^{\alpha (k-h)D} \left( N - h \right) e^{-\alpha D} \cos \left( 2b \frac{\omega}{c} D \right) \right],
\]

(1)

\[
Z_{IM}(\omega) = Z_0 \frac{\omega (\alpha_m + \alpha_e)}{4\pi^2 b^2 c} N \left[ \frac{1}{N} \frac{\omega}{4\pi b^2 \ln (d/b)c} - \frac{\omega}{(\alpha D)^2} \sum_{h=1}^{N-1} (N - h) e^{-\alpha D} \sin \left( 2b \frac{\omega}{c} D \right) \right],
\]

(2)

where \( D \) is the hole spacing and \( c \) is the speed of light. The attenuation constant \( \alpha \) accounts of any dissipation in the field propagation; in the case of ohmic losses at room temperature \( \alpha \) depends on the \( \sqrt{\omega} \). In practical cases, the ohmic dissipation is very small (the beam pipes are built with good conductors) and \( \alpha D \ll 1 \). Both real and imaginary parts depend on the interference among the holes leading to resonance peaks in the impedance at \( \omega_n = n\pi c/D \).

Far from the resonances, e.g., for \( \omega < \pi c/D \), the coupling impedance is well approximated by

\[
Z_{RE}(\omega) \simeq Z_0 \frac{\omega^2 (\alpha_m + \alpha_e)^2}{16\pi b^4 \ln (d/b)c^2} \left[ N + \frac{e^{-\alpha D N} - 1 + \alpha D N}{(\alpha D)^2} \right],
\]

(3)

\[
Z_{IM}(\omega) \simeq Z_0 \frac{\omega (\alpha_m + \alpha_e)}{4\pi^2 b^2 c} N,
\]

(4)

The real part grows up parabolically with the frequency; for small \( N \), such as \( \alpha D N \ll 1 \), the exponential term in Eq. (3) can be expanded up to the second order, resulting in an impedance proportional to \( N^2 \). Physically, this is due to the interference effect between the holes which dominates the real part of the impedance. On the contrary, the imaginary part is not affected (at least at a first order) by such an interference and it is mainly due to the holes as they were non interacting. Therefore the imaginary part grows up linearly with \( N \) and with the frequency \( \omega \), as in Eq. (4).

The loss factor, then, can be expressed as [3]

\[
k(\sigma) \simeq \frac{Z_0 \sqrt{\pi c}}{128\pi^4 b^4 \ln (d/b)\sigma} \left[ N^2 (\alpha_m + \alpha_e)^2 + \left( \frac{\sigma}{D} \right)^2 (\alpha_m - \alpha_e)^2 \right].
\]

(5)

Again, the dependence on \( N^2 \) takes into account interference effect between the holes and it is physically sound. Moreover, if

\[
\frac{1}{N^2} \left( \frac{\sigma}{D} \right)^2 \frac{(\alpha_m - \alpha_e)^2}{(\alpha_m + \alpha_e)^2} \ll 1,
\]

the loss factor is proportional to \( N^2 \), although it is not \( N^2 \) times the loss factor of a single hole (that is proportional to \( \alpha_m^2 + \alpha_e^2 \)). In formulae

\[
k(\sigma) \approx \frac{Z_0 \sqrt{\pi c}}{128\pi^4 b^4 \ln (d/b)\sigma} N^2 (\alpha_m + \alpha_e)^2
\]

(6)

if

\[
N \gg A \frac{\sigma}{D} \quad A = \frac{(\alpha_m - \alpha_e)}{(\alpha_m + \alpha_e)}
\]

In the case of circular holes in a thin wall, for example, \( A = 3 \).

The energy lost by the bunch in a section of length \( L_d = ND \) is the loss factor times \( Q^2 \), where \( Q \) is the charge of the bunch. Dividing that energy by the time the bunch takes to cover the length of the slotted pipe \( L_d \) (i.e., \( L_d/c \) for relativistic bunches) yields the power flowing in the coaxial and takes into account the effect of a single bunch traveling in the inner pipe. If we assume that the total power is the sum of the power due to each of the \( n_b \) bunches in a length \( L_d \), we get the expression:

\[
P = n_b \frac{Q^2 k(\sigma)}{L_d} = \frac{Q^2 k(\sigma)}{S_b},
\]

(7)

where \( S_b \) is the bunch spacing (\( S_b = L_d/n_b \)). In the limit of no attenuation and using Eq. (6) for the loss factor, the power is

\[
P = \frac{\sqrt{\pi}}{128\pi^4} \frac{Z_0 Q^2 c^2}{b^4 \ln (d/b)} \frac{(\alpha_m + \alpha_e)^2}{N^2}.
\]

(8)

**COLD TRANSITION NUMERICAL SIMULATIONS**

Quantitative results of the energy losses, parasitic resonances, longitudinal and transverse coupling impedance have been obtained by MAFIA simulations in time domain. Even if the code allows the study of relatively complicated 3D structures, to get accurate results long computation time and heavy memory resources are needed.

The simulations are not straightforward and require the choice of adequate mesh sizes. Numerical simulations of the cold transition were performed with fine mesh sizes (of the order of 2 mm) and considering only one half of the structure. Each simulation took about 30 hours of CPU time.

To evaluate the short range wake potential over the bunch length, in time domain simulations, a single Gaussian bunch with \( \sigma = 12 \) cm has been considered. The long range wake potential has been calculated by assuming a smaller bunch length (\( \sigma = 1 \) cm) over a distance of 6 m behind the bunch. The impedance of the structure has been estimated by the Fourier transform of this long range wake potential.

The dimensions of the SPS cold transition are reported in Fig. 2. It is a coaxial line with an inner elliptical cross section and with pumping holes located in the horizontal
plane. There are two rows of 73 holes with a diameter of 7 mm and with a distance among centers of 30 mm. The outer diameter of the coaxial line is about 100 mm, and the total length of the structure is 2232 mm. Since the structure needs a very large number of mesh points, to get a good accuracy and to better understand the results, we have considered four different cases assuming hole diameters equal to 8 mm:

- 18 holes (half structure) with 2 mm mesh sizes;
- 38 holes (half structure) with 2 mm mesh sizes;
- 73 holes (half structure) with 2 mm mesh sizes;
- 146 holes (half structure) with 4 mm longitudinal and 2 mm transverse mesh sizes.

We have also compared the numerical results with the analytical ones.

**Cold transition with 38 holes (19 holes per row)**

We have evaluated the power losses as a function of the holes number and as a function of the beam position to find if there is a maximum or minimum of these losses and to investigate if there is an interference effect among the holes.

In order to have a deeper insight of the involved physics, intensive calculations of the loss parameters as a function of the vertical and horizontal coordinates, of the beam positions and of the holes number have been carried out.

The longitudinal and transverse (vertical and horizontal) loss parameters as a function of the vertical and horizontal coordinates are reported in Figs. 4a, 4b, and 4c, considering a bunch length of 12 cm. Fig. 4a shows the longitudinal and transverse loss parameters (normalized to their maximum values) along the vertical coordinate of the symmetry plane. We can observe that the maximum value of the longitudinal loss factor occurs at the geometric center of the structure as it has to be expected because the distance between the beam and the holes is the smallest one. About the transverse loss factor we note that it is directly proportional to the derivative of the longitudinal one as the Panofsky - Wenzel theorem states for coupling impedances. In Figs. 4b and 4c the longitudinal and transverse loss factors as function of the horizontal coordinate (where the holes are located), are reported. It is very interesting to note that the longitudinal loss factor increases strongly (about two orders of magnitude) by going from the geometric center to the holes region while it is almost linear in the center of the pipe. This effect can cause heating of the transition during the machine operation if the orbit is not well corrected.

Calculating the average power \( P \) lost by the beam with the relation:

\[
P = K_l I^2 T N_b
\]

with a longitudinal loss parameter calculated at the geometric center \( K_l = 1.56 \times 10^4 \) V/C, an average current 0.7 mA per bunch, a revolution time \( T = 23 \times 10^{-6} \) s, and with a number of bunches \( N_b = 288 \), we get a power loss \( P = 0.051 \) mW. The gradient of the longitudinal loss factor with respect to the horizontal coordinate and near the center of the pipe is \( \partial K_l / \partial x = 2810 \) V/(C mm). As a consequence the power gradient is \( \partial P / \partial x = 9.1 \) \( \mu \)W/mm.

Concerning the transverse loss factor, close to the geometric center of the pipe, the horizontal one is three times higher than the vertical one. For this reason we present results related only to the horizontal case.

In the region where the horizontal loss factor exhibits a linear behavior, the following relation can be applied [5]:

\[
I_m(Z_t) = -2 \frac{\sigma K_t \sqrt{\pi}}{c}
\]

(10)
giving a transverse impedance of

\[
Z_t = 25.13 \text{ } \Omega/m,
\]

with \( \partial K_t / \partial x = 1.77 \times 10^{10} \text{ } \text{V/}(\text{Cm}) \)

As far as the longitudinal coupling impedance is concerned, we calculated the short range wake potential over the bunch length by assuming \( \sigma = 12 \) cm.

The normalized wake potential is reported in Fig. 4d. It is clear that the transition is purely inductive at low frequencies.

From the simulations we have that the absolute maximum values of the wake potential are \( |W_{\text{max}}| = |W_{\text{min}}| = W = 1.831 \times 10^7 \text{ V/C} \). By using the theoretical relation [4]:

\[
W = \frac{Lc^2}{\sqrt{2\pi\epsilon\sigma^2}}
\]

(11)
we can estimate the inductive longitudinal impedance \( |Z/n| = \omega_0 L = 3.31 \times 10^{-6} \text{ } \Omega \).

**Cold transition with 74 holes (37 per row)**

To investigate the behavior of the longitudinal and transverse loss factors as a function of the holes number, simulations were performed with 74 holes. The corresponding results are still reported in Figs. 4a, 4b, 4c, and 4d.

Fig. 4a shows the longitudinal and transverse loss parameters normalized to their maximum values on the vertical symmetry plane. These results reproduce the behavior of the previous case with 38 holes. The calculated longitudinal loss factor in the center of the pipe is \( K_l = 5.132 \times 10^4 \) V/C and it yields a power losses \( P = 0.17 \) mW. The gradient of \( K_l \) (see Fig. 4b) as a function of the horizontal coordinate is \( \partial K_l / \partial x = 9229 \) V/(C mm) and the related the power gradient is \( \partial P / \partial x = 30 \) \( \mu \)W/mm. The calculated transverse impedance (see Fig. 4c) is \( Z_t = 49 \text{ } \Omega/m \) with \( \partial K_t / \partial x = 3.45 \times 10^{10} \text{ V/}(\text{Cm}) \), while the longitudinal impedance is \( |Z/n| = 6.4 \times 10^{-6} \text{ V/C} \) with \( |W_{\text{max}}| = |W_{\text{min}}| = 3.565 \times 10^7 \text{ V/C} \).

The coupling impedance at low frequencies is shown in Fig. 5 for 3 different values of the hole number \( N \). The parabolic behavior is clear from the plots and can be highlighted by a polynomial fit. The result is shown in Table 1 where the correlation coefficient \( R \) states the quality of the fit: it is equal to 1 for an optimal fit.
Parasitic resonances

The ratio $a_i/a_j$ is very close to $(N_i/N_j)^2$ for $i, j = 1, 2, 3$, as it should be.

Parasitic resonances

To completely characterize the structure, we have calculated the 6 m long range wake potential considering a Gaussian bunch of $\sigma = 1$ cm traveling through the structure 4 mm away from the holes; in fact, such a short distance allows a better measurement of the resonance amplitudes. The Fourier transform of this wake potential gives the impedance of the structure at that beam position. The cutoff frequencies of the vacuum chamber calculated by OSCAR2D [6] are 2.11 GHz for the TE mode polarized with the E-field parallel to the major axis and 2.63 GHz for the other TE mode.

The real and imaginary part of the longitudinal coupling impedance and the normalized wake potential with 38 holes are illustrated in Figs. 6a, 6b, 6c. The results obtained with 74 holes are shown in Figs. 7a, 7b, 7c. The results obtained with 74 holes are shown in Figs. 8a, 8b, 8c for the longitudinal case and in Figs. 9a, 9b, 9c for the transverse one. Comparing the results of the pictures related to 38 and 74 holes, one observes that at low frequency (up to $\sim 4$ GHz) the impedance is purely inductive. At low frequency the imaginary part scales with $\omega$, as expected from theory. The real part, as already observed, is proportional to $\omega^2$ and exhibits resonant peaks. The maximum amplitude of the resonances depends on the proximity.
of the beam to the slotted wall, but they are always present. In the simulations, their amplitude scales with $N$, due to the small number of points used to resolve the resonances.

The plots clearly show some parasitic resonances in the 4 - 12 GHz frequency range very far from the bunch spectrum cut-off. It is worth noticing that the strongest resonance is peaked at about $f = 9.27$ GHz. This frequency corresponds to a wavelength equal to the holes distance. For an additional confirmation of this parasitic resonance related to the holes distance, we have simulated the case of a 28 mm distance between the holes. As a result, the corresponding frequencies of all the parasitic resonances scale proportionally to the distance between two hole centers.

The impedances follow the same dependence with $N$ as discussed in the previous section.

Moreover, Fig. 10 shows the dependence of the term $\alpha ND$ in the frequency range of the SPS bunch spectrum for some values of the hole number. In fact, even at a frequency $f = 2$ GHz and at room temperature, we get an attenuation constant $\alpha = 4.8 \times 10^{-3} \text{ m}^{-1}$ which satisfies the theoretical condition $\alpha ND << 1$.

For increasing our confidence with the presented analysis, we have simulated a coaxial outer short circuited with a length of $l = 1.3$ m. The corresponding results are illustrated in Figs. 11a, 11b, 11c, and 12a, 12b, 12c.

As predictable, it is possible to observe a modulation of
the impedance \( c/(2l) = 115 \text{ MHz} \) and higher quality factors of the parasitic modes because the resonant fields do not propagate in the coaxial line.

**Whole cold transition with 146 holes (73 holes per row)**

For the whole structure, a 4 mm longitudinal mesh was used due to limitation of the number of mesh points. Fig. 13a shows the longitudinal loss parameter and the power losses as a function of the horizontal displacement. It is clear that, for an off-set beam, the power losses increase considerably as it has been already discussed. As an example, a beam with \( x = 2.4 \text{ cm} \) off-set loses about \( P = 3.2 \text{ mW} \) that can affect machine operation.

To investigate the behavior of the power losses as a function of the beam pipe shape, different simulations were done changing the smaller axis of the elliptical beam pipe. Fig. 13b shows the power losses as function of the horizontal displacement for different values of the smaller axis \( a \). The case with \( a = 36 \text{ mm} \) corresponds to the actual transition while that one with \( a = 42 \text{ mm} \) to a circular cross section.

We can observe that the circular shape gives less power losses than the actual one. This is an important element for the final section shape choice. We can observe that, in the elliptical cross section and in the center of the beam pipe, the power loss is \( P \sim 0.4 \text{ mW} \) (0.16 mW/m on overall transition) and the power gradient along the horizontal co-
ordinate $\partial P/\partial x = 0.054 \text{ mW/mm}$ with $\partial K_t/\partial x = 1.65 \times 10^4 \text{ V/(Cmm)}$. Again, these values scale approximately as the square number of holes $N^2$.

Fig. 13c shows the inductive behavior of the wake potential with $\sigma = 12 \text{ cm}$. The spikes observed are due to the large longitudinal mesh adopted for the present case. By assuming $|W_{\text{max}}| = |W_{\text{min}}| = 7.3 \times 10^7 \text{ V/C}$, we obtain $|Z/\sigma| = 13 \times 10^{-6} \text{ } \Omega$ which again scales as the hole number $N$. This result is also confirmed by the Fourier transform of the wake potential obtained with $\sigma = 2 \text{ cm}$ over a distance of few meters behind the bunch which gives $|Z/\sigma| = 12 \times 10^{-6} \text{ } \Omega$ (Fig. 13d).

In Fig. 14 we have reported the transverse loss parameter as function of the transverse displacement. The horizontal impedance is $Z_{tx} = 99 \text{ } \Omega/m$ with $\partial K_t/\partial x = 7 \times 10^{10} \text{ V/(Cm)}$ while the vertical one is $Z_{ty} = 30 \text{ } \Omega/m$ with $\partial K_t/\partial x = 2 \times 10^{10} \text{ V/(Cm)}$. These values scale like the hole number $N$. A check of these results by using the Fourier transform of the wake potential gave $100 \text{ } \Omega/m$ and $29 \text{ } \Omega/m$, respectively.

**KOLD TRANSITION ANALYTICAL ESTIMATIONS**

Even though in the present study the chamber has a beam pipe with elliptical cross-section, and therefore a correct approach to the solution of the problem needs three-dimensional codes, the analytical methods remain extremely useful to check the order of magnitude of the discussed numerical results.

We remark that the analytical model presented in section 2 is valid for a coaxial beam pipe with axial symmetry, it neglects the effects due to the curvature of the chamber ($a << b$ being $a$ the hole radius and $b$ the chamber radius),

---

Figure 11: Longitudinal coupling impedance for $N = 74$ with a coaxial length $l = 1.3 \text{ m}$: a) Real part, b) Imaginary part, c) Scaled wake potential.
it assumes zero wall thickness of the circular hole and finally it considers the approximation of an incident plane wave (fields constant over the hole).

Nevertheless, for the longitudinal analysis, if we introduce an equivalent circular radius due to the ellipticity of the beam pipe [7], and by taking into account the wall thickness correction factors on the polarizabilities [8], we can still apply the theoretical model.

As a result for the cold transition with a number of holes $N = 146$, by assuming an equivalent chamber radius $b = b_{eq} = 0.0374$ m and by keeping unchanged the outer one (Fig. 2), with a hole wall thickness of 2 mm, we get $Z/n = 16.5$ $\mu$\Omega and a power loss $P = 0.43$ mW. By assuming instead an internal radius equal to $b = 0.043$ m, we obtain $Z/n = 12.3$ $\mu$\Omega and $P = 0.55$ mW. The corresponding numerical estimations (about $Z/n = 12$ $\mu$\Omega and $P = 0.4$ mW) are in good agreement with the previous ones.

Figure 12: Transverse coupling impedance for $N = 74$ with a coaxial length $l = 1.3$ m: a) Real part, b) Imaginary part, c) Scaled wake potential.

Figure 13: Whole cold transition with $N = 146$: a) Longitudinal loss parameter and power losses as a function of the horizontal displacement, b) Power losses as a function of the horizontal displacement for some values of the vertical axis by keeping unchanged the horizontal one, c) Wake potential normalized to its maximum value as a function of the distance from bunch head ($W_{max} = 7.3 \times 10^5$ V/C), d) Fit of the imaginary part of the longitudinal broad band coupling impedance.
CONCLUSION

We presented the study of the cold to warm transitions of the SPS machine. The numerical estimations of the coupling impedance have been compared to a theoretical model showing a good agreement. Large beam off-set deposits higher power. The obtained ohmic losses are much less than those measured in the SPS cold transitions. However even losses due to the pumping holes (of the order of few tens of mW per meter) lead to heat overload. Therefore it is advisable an optimization of the beam pipe shape.

REFERENCES

Coherent Beam-Beam Modes in the LHC for Multiple Bunches, Different Collision Schemes and Machine Symmetries

W. Herr and T. Pieloni, CERN, CH-1211 Geneva (CH)

Abstract

In the LHC almost 3000 bunches in each beam will collide near several experimental regions and experience head-on as well as long range beam-beam interactions. In addition to single bunch phenomena, coherent bunch oscillations can be excited. Due to the irregular filling pattern and the asymmetric collision scheme, a large number of possible modes must be expected, with possible consequences for beam measurements. To study these effects, a simulation program was developed which allows to evaluate the interaction of many bunches. It is flexible enough to easily implement any possible bunch configuration and collision schedule and also to study the effect of machine imperfections such as optical asymmetries. First results will be presented and future developments are discussed.

INTRODUCTION

The spectra of the barycentric motion and the mode frequencies of coherent beam-beam modes are well known and understood for the case of a few bunches colliding head-on [1, 2]. Present and future colliders have many bunches and multiple interaction points and a much richer spectrum of modes must be expected [3]. This is in particular true when the collision points are not symmetrically distributed and additional effects due to non-symmetric collision schemes [4, 5] or asymmetric colliders like two-ring schemes [6] must be expected.

In the LHC there are a number of effects which break the symmetry between the collision points:

- Asymmetric configuration of the collision points
- Presence of a large number of parasitic long range interactions
- Unavoidable PACMAN effects [7, 8]
- It is impossible to make the bunches collide exactly head-on [9, 10].

In the case of multiple head-on collisions these modes can be analyzed with a linearized model searching for the eigenmodes of the full single turn map. However, when the non-linear long range interactions are included, the linearized treatment is not adequate. One therefore might expect a fairly large number of modes which may obscure tune measurements or feedback systems. The presence of a large number of modes due to the effect of local, parasitic interactions was already studied in [11, 12] but without possible PACMAN effects and for a simplified LHC collision scheme. It is therefore important to define possible configurations which minimize the number of modes and provide cleaner spectra. For the evaluation a strong-strong simulation program is written, using a rigid Gaussian model for the bunches.

SIMULATION PROGRAM

The simulation program must allow to:

- Track each bunch of both beams independently around the ring
- Apply head-on and long range interactions at bunch encounters
- Give initial kicks to single bunches or a range of bunches to simulate excitation (e.g., for tune measurement)
- Analyze the motion of selected or a range of bunches
- Show the complete set of possible coupled beam-beam modes

In order to evaluate different scenarios, the program must be very flexible to allow easy changes of parameters such as tunes, number of bunches, filling scheme, collision scheme etc. In particular it must allow different crossing planes. The possibility to change the phase advance between collision points is important. Statistical fluctuations such as bunch intensity, emittance etc. must be possible to simulate. It should be possible to simulate and demonstrate PACMAN effects. In order to get all correct modes of the bunches coupled by head-on and long range interactions, all individual interactions must be simulated in full. In particular, lumping several long range interactions is therefore not adequate. For future extensions it must be possible to add multi-particles to replace rigid bunches in a straightforward way.

Parameters

To describe the motion of a rigid bunch the following parameters are used:

- Horizontal position and angle of barycentre: X and X'
- Vertical position and angle of barycentre: Y and Y'
- Horizontal position and angle of single particles: x and x'
• Vertical position and angle of single particles: y and \( y' \)

• Longitudinal phase (or position \( s \)) and energy deviation: \( \phi \) (or \( s \)) and \( \delta \)

For extension and later use, following parameters are foreseen and stored:

• Bunch intensity (to determine beam-beam kick)
• Bunch emittance (to determine beam-beam kick)
• Tune shift \( \Delta Q_X \) and \( \Delta Q_Y \) with respect to a nominal bunch.

**Input description**

For the simulation it is necessary to describe the arrangement of the bunches around the machine and their possible interactions with other bunches or machine elements. For simplicity it must be optimized to study beam-beam interactions. However, the description should be very flexible to allow the study of different filling or collision schemes as well as optical properties of the machine. I have followed the strategy designed for beam-beam tracking and the computation of self-consistent properties [9, 10, 13, 14] and included all the necessary extensions.

**Bunches in the ring and description of filling scheme**

The description of the bunch filling scheme is given in the form of groups. Each group has two parameters: the first specifies the number of slots \( n \) and the second whether the \( n \) slots are occupied by a bunch (1) or whether the slots are empty (0). The total number of slots must be equal to the machine circumference divided by the bunch spacing. It is therefore vital that all empty slots are defined as well as all filled slots. The number of groups per line is specified at the beginning of the description file. To define 1 bunch followed by 39 empty slots one could use:

```plaintext
# bunch filling example 1
# Number of groups
2
1 1 39 0
1 1 39 0
1 1 39 0
```

This example describes 4 equidistant bunches spread out in 160 slots (possible bunch positions) while the example below represents the actual LHC bunch filling scheme [7, 18, 19].

```plaintext
# bunch LHC filling example
# number of groups
8
72 0 8 0 72 1 8 0 72 1 8 0 30 0 0 0
```

The number of slots in this case is 3564 which is one tenth of the LHC harmonic number. The description should maximize the readability, although any format is possible.

**Positions and actions**

When one is interested in beam-beam interactions, only every half bunch spacing something can happen (i.e. where two bunches from the two beams could meet). For \( N \) slots defined by the filling scheme (i.e. number of possible bunch positions), one has \( 2N \) positions where such actions can occur. In the description the numbering of the positions goes from 1 to 2\( N \) in the direction of the clock-wise beam.

**Definition of actions**

At any position, an action can be requested for a bunch when it is in that place. For beam-beam interactions (head-on or long range) two bunches (i.e., one from each beam) must be at this position. The different actions are specified by a code number. Possible actions are:

- Head-on collision (at the specified position, code 2 or -2)
- Head-on and long range collisions (left and right of a specified head-on collision)
- Multiple long range collisions (left and right of a specified position, code 4 or -4)
- Single separated collision (code 5 or -5)
- Linear matrix transfer of a bunch (code 3)
- No action (default)

Additional actions, e.g., non-linear elements or correction devices, can easily be defined.

**Head-on collision**

The code for a head-on collision point is either 2 or -2. The positive sign indicates horizontal and the negative sign vertical separation of the associated long range interactions, i.e., crossing plane in the case of the LHC. The strength of the head-on collision is determined by the beam-beam parameters which is either taken from the general input file.
calculated from the bunch intensities, emittances and positions of the two colliding bunches. Before and after a head-on collision, the bunches are advanced in transverse phase space by $\pi/2$.

**Long range collisions left and right of a head-on collision:**

When a head-on collision point is defined like above, a number of long range collisions left and right of the collision point can be specified on the action statement for the head-on collision by specifying the number of collision points, i.e., the number of positions where long range interactions can occur, e.g., the line:

```
161 2 -15 +15
```

specifies a head-on collision at position 161 with horizontal crossing and 15 long range interactions on each side.

**Long range collisions left and right of a specified position:**

When an action code of 4 or 4 is specified, only the long range interactions left and right of a specified position are active, the central head-on collision is ignored. This can be used to simulate a crossing angle configuration when the central head-on collision point is separated and the bunches experience long range interactions left and right of the symmetry point. However the rotation by $\pi/2$ before and after the specified position is performed to ensure the correct phase relationship between the long range interactions before and after.

**Separated collisions:**

An action code of 5 or 5 is used for a single separated interaction (e.g., in a Pretzel scheme). The third and fourth parameters are ignored.

**Linear transfer of the bunches:**

With the action code 3 a linear transfer is defined. The two parameters are used to control the phase advance of the transfer. The parameters specify the phase advance in units of $2\pi$ (tune). The phase advance between any point in the machine and in particular between interaction points can easily be controlled that way. The two rotations of $\pi/2$ for each head-on interaction point must be taken into account to get the correct overall tune. In the present implementation the phase advance between two points in the machine is assumed to be the same for the forward and backward beams. In a two ring machine like the LHC this is not always the same.

**Description of collision scheme**

The collision scheme defines the actions to be performed at the possible positions. This description is an extension of the scheme defined for [13]. Every action consists of one line which defines first the position of the action, the second column is the code of the desired action and the third and fourth columns are parameters required by the action. Typical collision descriptions are:

### Collision scheme 1 (for filling example 1):

```
1 2 -5 +5
21 3 7.535 6.91375
41 -2 -5 +5
101 3 23.605 21.74125
161 -2 -5 +5
221 3 23.605 21.74125
281 2 -0 +0
301 3 7.535 6.91375
```

which defines 4 collision points where three have long range collisions on both sides of the head-on collision points. The machine has an eightfold symmetry in geometry and phase advance. As a further example, a collision scheme representing the LHC with its present filling scheme and layout of the four experiments is shown below.

### Collision scheme LHC (for LHC filling scheme):

```
1 -2 -15 +15
447 3 8.046 6.940
882 -2 -0 +0
2229 3 23.015 21.821
3565 2 -15 +15
4902 3 23.533 20.689
6236 2 -0 +0
6684 3 7.716 7.870
```

Since the filling scheme defines the number of bunches and positions, the collision definition scheme must always follow the definition of the filling pattern.

**Parameter input**

At the start of the program, a parameter file is read in to define the basic input data. The name of this file is taken as a command line argument of the program,

```
collision: coll_ref.in  //input collision
filling: fill_ref.in  //beam filling
use bunch: 1  //bunch to analyze
number of turns: 14  //2**14 turns
bunches to kick: 5  //5 bunches kicked
sigma int: 0.8  //intensity fluctuations
b-b parameter: 0.0025  //beam-beam parameter
```

**Actions**

**Linear transfer**

At a position requiring a linear transfer we use a linear transfer map. If we define for example A as the linear transfer matrix for the horizontal coordinates where $\Delta \mu_X$ is taken from the input files (i.e., collision scheme):
\[ A = \begin{pmatrix} \cos(\Delta\mu x) & \sin(\Delta\mu x) \\ -\sin(\Delta\mu x) & \cos(\Delta\mu x) \end{pmatrix} \]  

Then we obtain for the horizontal coordinates the following expression:

\[ \begin{pmatrix} X \\ X' \end{pmatrix}_{n+1} = A \begin{pmatrix} X \\ X' \end{pmatrix}_n \]  

The same transfer map with \( \Delta\mu y \) is applied to the vertical coordinates \( Y \) and \( Y' \).

**Head-on beam-beam interaction**

To calculate the head-on beam-beam kick on a bunch, the counter-rotating beam distribution is assumed to have a Gaussian density distribution in the two planes with barycentres at \( (X^*, Y^*) \) and squared transverse sizes \( \Sigma_{xx} = \langle (x - X)^2 \rangle \) and \( \Sigma_{yy} = \langle (y - Y)^2 \rangle \). In that case the beam-beam force can be expressed analytically. The * denotes parameters of the opposing beam. In the case of rigid bunches the transverse sizes are kept constant. We apply a horizontal beam-beam kick at the IP (equivalent for the vertical beam-beam kick):

\[ \Delta X' = \frac{2r_p N_p^x \beta_x}{\gamma \sigma_x^2} F_x(X - X^*, Y - Y^*, \Sigma_{xx}, \Sigma_{yy}) \]  

with \( r_p \) the classical proton radius, \( N_p^x \) the bunch population (* indicates parameters of the counter-rotating beam), \( \gamma \) is the relativistic Lorentz factor, \( \beta_x \) the horizontal beta function at the IP, \( \sigma_x \) the horizontal rms size and \( F_x \) (or equivalently, \( F_y \) for the vertical beam-beam kick) given by

\[ F_{(x,y)}(X - X^*, Y - Y^*, \Sigma_{xx}, \Sigma_{yy}) = \frac{(X, Y)}{(X^2 + Y^2)} \left[ 1 - \exp \left( -\frac{X^2 + Y^2}{\Sigma_{xx}^2 + \Sigma_{yy}^2} \right) \right] \]  

which is the expression for round beams when \( \Sigma_{xx} \approx \Sigma_{yy} \). When the beams are not round, we use the Bassetti-Erskine formula for the evaluation of the kick [22]. In the horizontal plane. The map at the beam-beam interaction is then:

\[ \begin{pmatrix} X \\ X' \\ Y \\ Y' \end{pmatrix}_n = \begin{pmatrix} X \\ X' + \Delta X' \\ Y \\ Y' + \Delta Y' \end{pmatrix}_n \]  

The beam-beam parameters are defined by

\[ \xi_{(x,y)} = \frac{N_p^x \beta_x \beta_{(x,y)}}{2 \pi \gamma \sigma_{(x,y)} (\sigma_x + \sigma_y)} \]  

With the nominal LHC parameters we have \( \xi \approx 0.0034 \).

Before and after each head-on collision, I apply a phase advance of \( \pi/2 \) in each plane.

**Long range beam-beam interaction**

For the calculation of the long range beam-beam kick, the expressions for the head-on interaction must be modified to take the separation into account. The constant part of the kick in the plane of separation must be subtracted. Assuming a constant horizontal separation \( d \) and using the expression (4):

\[ \Delta X' = \frac{2r_p N_p^x \beta_x}{\gamma \sigma_x^2} [F_x(X + d - X^*, Y - Y^*, \Sigma_{xx}, \Sigma_{yy}) - F_x(d, 0, \Sigma_{xx}, \Sigma_{yy})] \]  

one gets the deflection \( \Delta X' \) and for the other plane we have:

\[ \Delta Y' = \frac{2r_p N_p^y \beta_y}{\gamma \sigma_y^2} [F_y(X + d - X^*, Y - Y^*, \Sigma_{xx}, \Sigma_{yy})] \]  

**Tracking strategies**

**Initial conditions**

The barycentres of the bunches of the two beams can be set all to zero at the start of the program or distributed according to a Gaussian distribution. Furthermore, a single bunch or a small number of bunches can be excited at the beginning, simulating e.g. a tune measurement. This is specified in the input file.

**Rotation of bunches in both rings**

Beam 1 bunches travel increasing number of positions, beam 2 bunches decreasing number of positions. At each step, every bunch is advanced by one position, i.e. half a bunch spacing. One complete turn in the machine therefore requires 2N steps. The calculations for all bunches of a beam at each step are independent and it can be envisaged to make use of parallel processing, in particular when the bunches consist of many macro-particles in a later version of the program.

**Data processing**

By a Fourier analysis of the barycentre of the bunches, as calculated turn by turn, we obtain the tune spectra of the dipole modes. For only one bunch per beam the two spectra of the two bunches are equivalent. For more than one bunch per beam the spectra of bunches with the same collision scheme are also equivalent. Analyzing the sum \( (X^{(1)} + X^{(2)}) \) or the difference \( (X^{(1)} - X^{(2)}) \) of the barycentre of two colliding bunches of the two beams (denoted by (1) and (2)) show the spectra of the 0- and \( \pi \)-modes separately. This is useful to analyze the details of the modes.
Program validation

Head-on effects

To validate the program we simulated the simple and well known case of two head-on collisions in two interaction points, opposite in azimuth. A value of 0.0025 was used for the linear beam-beam parameter. The result of the simulation is shown in Fig. 1. Equal charges of the two colliding beams are assumed and the frequencies are therefore shifted downwards from the unperturbed tunes. In Fig. 1 the spectrum of the first bunch of beam

![Figure 1: Symmetric Head-on collisions in IPs 1 and 5.](image1)

1 is shown; the spectra clearly shows the two coherent beam-beam modes. The sum signal of the two bunches gives the so called 0-mode (right peak in Fig. 1) while the difference signal gives only the π-mode signal (left peak in Fig. 1). This is in agreement with the expectations. The frequency shift between the 0-mode and the π-mode is however not correct in a rigid bunch model and the forces must be calculated from the real field distribution [16, 17]. For the purpose of this report to study the spectra of dipole oscillations the rigid Gaussian model is adequate.

RESULTS

With the available simulation program, the following effects can be studied and the dependence of the results on the optical and collision configuration can be evaluated:

- Head-on interactions only (one bunch per train or no long range positions)
- Head-on and long range interactions (multiple bunches per train)
- Excitation of single and multiple bunches in a train for measurement purposes

Multiple head-on interactions

In the case of multiple head-on collisions in a machine, the symmetry properties of the layout are very important for the spectra. A high degree of symmetry can lead to the degeneracy of modes, i.e. identical frequencies, and their suppression in the spectra. Breaking the symmetry by choosing a non-symmetric collision scheme or phase advance differences between the interaction points may cancel this effect and leads to the appearance of additional modes in the spectra. In the following I assume a collider with an eightfold symmetry of the possible collision points and number the interaction regions from 1 to 8. In this case the interaction points 1 and 5 are opposite in azimuth. This resembles the geometrical layout of the LHC straight sections. The Fig. 1 shows two head-on collisions opposite in azimuth with symmetric optical layout. The

![Figure 2: Non symmetric head-on in IPs 1 and 2.](image2)

![Figure 3: Symmetric head-on in IPs 1, 3, 5 and 7.](image3)

effect of additional head-on collisions and non symmetries in the optical layout are shown in the horizontal spectra in Figs. 2 and 3 are shown for two simple cases. It can be observed that a higher degree of symmetry (or periodicity) leads to degeneracy of mode frequencies and fewer spectral lines. This confirms earlier findings [4, 5, 11, 12] and the importance of symmetries for coherent modes. The number of lines in the spectra can be qualitatively under-
stood by analysing the collision pattern of the bunches. The number is closely related to the number of bunches to which the measured bunch couples directly or indirectly (i.e. via other bunches). For example this explains the number of spectral lines when collisions occur only in interaction points 1 and 2 (Fig.2) and the reduced number when collisions occur in points 1, 3, 5 and 7 (Fig.2).

**Collisions with the LHC interaction region layout**

The collision scheme of the LHC with its four interaction regions was illustrated as an example already. Although the geometry has an eightfold symmetry, the phase advances between the interaction points break this symmetry and we must expect a richer spectrum of modes.

**LHC interaction region layout with standard phase advance**

The standard LHC collision scheme was already used as an example before. For the tracking studies, the arcs can be compressed since no action can happen except a single linear transfer. The number of bunches is reduced to 9 per train and to observe PACMAN effects, the number of long range positions is 5 on each side of the collision point. This will strongly reduce the required computing time but has no qualitative effect on the results. The nominal collision definition scheme used in the simulation is then:

<table>
<thead>
<tr>
<th></th>
<th>1</th>
<th>2</th>
<th>5</th>
<th>+5</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>2</td>
<td>-5</td>
<td></td>
<td></td>
</tr>
<tr>
<td>41</td>
<td>3</td>
<td>8.046</td>
<td>6.940</td>
<td></td>
</tr>
<tr>
<td>81</td>
<td>-2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>202</td>
<td>3</td>
<td>23.015</td>
<td>21.821</td>
<td></td>
</tr>
<tr>
<td>321</td>
<td>-2</td>
<td>-5</td>
<td>+5</td>
<td></td>
</tr>
<tr>
<td>441</td>
<td>3</td>
<td>23.533</td>
<td>20.689</td>
<td></td>
</tr>
<tr>
<td>561</td>
<td>2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>601</td>
<td>3</td>
<td>7.716</td>
<td>7.870</td>
<td></td>
</tr>
</tbody>
</table>

together with a filling scheme:

<table>
<thead>
<tr>
<th>Number of groups</th>
</tr>
</thead>
<tbody>
<tr>
<td>2</td>
</tr>
<tr>
<td>9 1 71 0</td>
</tr>
<tr>
<td>9 1 71 0</td>
</tr>
<tr>
<td>9 1 71 0</td>
</tr>
<tr>
<td>9 1 71 0</td>
</tr>
</tbody>
</table>

**LHC interaction region layout with symmetry between IP1 and IP5**

Starting from the scheme above, it can be partially symmetrized to fulfill:

$$\Delta Q_x^{+5} = \Delta Q_x^{-5} = Q_x / 2$$

and we use:

<table>
<thead>
<tr>
<th></th>
<th>1</th>
<th>-2</th>
<th>-0</th>
<th>+0</th>
</tr>
</thead>
<tbody>
<tr>
<td>41</td>
<td>3</td>
<td>8.046</td>
<td>6.940</td>
<td></td>
</tr>
</tbody>
</table>

**Figure 4:** Head-on collisions in IPs 1, 2, 5 and 8 with nominal LHC phase advance between interaction points.

<table>
<thead>
<tr>
<th></th>
<th>81</th>
<th>-2</th>
<th>-0</th>
<th>+0</th>
</tr>
</thead>
<tbody>
<tr>
<td>201</td>
<td>3</td>
<td>23.109</td>
<td>21.720</td>
<td></td>
</tr>
<tr>
<td>321</td>
<td>2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>441</td>
<td>3</td>
<td>23.439</td>
<td>20.790</td>
<td></td>
</tr>
<tr>
<td>561</td>
<td>2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>601</td>
<td>3</td>
<td>7.716</td>
<td>7.870</td>
<td></td>
</tr>
</tbody>
</table>

A further improvement is possible by adjusting the phase advance between interaction points 2 and 8 as shown below:

<table>
<thead>
<tr>
<th></th>
<th>1</th>
<th>-2</th>
<th>-0</th>
<th>+0</th>
</tr>
</thead>
<tbody>
<tr>
<td>41</td>
<td>3</td>
<td>8.046</td>
<td>6.940</td>
<td></td>
</tr>
<tr>
<td>81</td>
<td>-2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>201</td>
<td>3</td>
<td>23.109</td>
<td>21.720</td>
<td></td>
</tr>
<tr>
<td>321</td>
<td>2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>441</td>
<td>3</td>
<td>23.6235</td>
<td>21.270</td>
<td></td>
</tr>
<tr>
<td>561</td>
<td>2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>601</td>
<td>3</td>
<td>7.5315</td>
<td>7.390</td>
<td></td>
</tr>
</tbody>
</table>

The spectra for such a scheme are shown in Fig.5. The comparison between Fig.4 and 5 shows the effect of the symmetry between interaction points 1 and 5. Although the number of modes is not really changed, it must be expected that the Landau clamping of modes with frequencies just below the 0-mode will "clean" the spectra around the 0-mode, i.e. the nominal tune, and therefore simplifies the tune measurements.

**LHC interaction region layout with full eightfold symmetry**

The fully symmetric version with eightfold symmetry in the phase advances is:

<table>
<thead>
<tr>
<th></th>
<th>1</th>
<th>-2</th>
<th>-0</th>
<th>+0</th>
</tr>
</thead>
<tbody>
<tr>
<td>41</td>
<td>3</td>
<td>7.78875</td>
<td>7.165</td>
<td></td>
</tr>
<tr>
<td>81</td>
<td>-2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
<tr>
<td>201</td>
<td>3</td>
<td>23.36625</td>
<td>21.495</td>
<td></td>
</tr>
<tr>
<td>321</td>
<td>2</td>
<td>-0</td>
<td>+0</td>
<td></td>
</tr>
</tbody>
</table>
Figure 5: Head-on collisions in IPs 1, 2, 5 and 8. Phase advance symmetry restored between IP1 and IP5 and adjusted between 2 and 8.

441  3  23,36625  21,495
561  2  7,78875  7,165

Figure 6: Head-on collisions in IPs 1, 2, 5 and 8 with full eight fold symmetry of phase advance.

The spectra for the fully symmetric machine are very similar to those obtained with the 'tuned' collision scheme shown in Fig.5.

SUMMARY

We have used a multi bunch simulation to compute the spectra of dipole oscillations driven by head-on and long range beam-beam interactions. The spectra largely depend on these interactions and the main observations can be summarized:

- Configuration of collisions should be symmetric to reduce number of dipole modes.
- Phase advance between low \( \beta \) interaction regions should be symmetric to allow degeneracy and compensation of coherent modes.
- Although not required for the compensation of first order PACMAN effects ([20]) or suppression of resonances ([5]), some flexibility of the phase adjustment between interaction points is desirable.
- Measurement should be on a single bunch following an excitation of this bunch if possible.

These should serve as recommendations when it becomes important to keep the spectra clean. However, damping effects such as Landau damping due to the incoherent tune spread are not included in the present rigid bunch model and will be studied in the future using a multi-bunch, multiparticle simulation. It must be expected that these damping effects suppress a significant number of modes, in particular in the immediate neighbourhood of the 0-mode.

REFERENCES

[12] W. Herr; Coherent dipole oscillations and orbit effects induced by long range beam-beam interactions in the LHC, CERN SL/91-34 (AP) and LHC Note 165 (1991).
HIGH BRIGHTNESS PROTON BEAMS FOR LHC: NEEDS AND MEANS

M. Benedikt, R. Garoby, CERN, Geneva, Switzerland

Abstract
Experiments [1, 2] have proven that the LHC injector chain can deliver a proton beam with the nominal characteristics (bunch intensity $N_b=1.15 \times 10^{11}$ protons per bunch (ppb) in normalised rms transverse emittances of $\varepsilon_{x,y}=3.5 \mu mrad$), but cannot reach the ultimate performance ($1.7 \times 10^{11}$ ppb in the same emittances). Moreover, in the longer term, an even higher beam brightness is required by all methods considered for increasing the LHC luminosity beyond the present ultimate level. Improvements and/or new processes are therefore needed, especially in the low energy accelerators. A number of solutions have already been imagined for the PS complex that involve new linac(s) or/and sophisticated beam gymnastics. The present capabilities and limitations of the accelerator chain are described. The needs of the possible LHC luminosity upgrades are outlined, the proposed improvements are explained and their features and performance are compared.

CAPABILITIES OF THE EXISTING ACCELERATOR CHAIN
The characteristics of the beam circulating in the LHC for nominal and ultimate luminosity ($1.0 \times 10^{34} \text{cm}^{-2}\text{s}^{-1}$ and $2.3 \times 10^{34} \text{cm}^{-2}\text{s}^{-1}$ respectively) are compiled in Table 1 [3]. To meet these requests, extensive modifications have been implemented in the complex of existing accelerators that serve as LHC injectors [3, 4].

Table 1: Nominal and ultimate proton beam characteristics in the LHC

<table>
<thead>
<tr>
<th>Injection</th>
<th>Collision</th>
</tr>
</thead>
<tbody>
<tr>
<td>GeV</td>
<td>GeV</td>
</tr>
<tr>
<td>Luminosity</td>
<td></td>
</tr>
<tr>
<td>nominal</td>
<td>450</td>
</tr>
<tr>
<td>ultimate</td>
<td>7000</td>
</tr>
<tr>
<td>Number of bunches</td>
<td>2808</td>
</tr>
<tr>
<td>Minimum bunch spacing</td>
<td>24.95 ns</td>
</tr>
<tr>
<td>$N_b$ (Intensity/bunch)</td>
<td></td>
</tr>
<tr>
<td>nominal</td>
<td>$1.15 \times 10^{11}$</td>
</tr>
<tr>
<td>ultimate</td>
<td>$1.7 \times 10^{11}$</td>
</tr>
<tr>
<td>Beam current</td>
<td></td>
</tr>
<tr>
<td>nominal</td>
<td>0.58</td>
</tr>
<tr>
<td>ultimate</td>
<td>0.86</td>
</tr>
<tr>
<td>$\varepsilon_x^*$ (normalised rms transverse emittances)</td>
<td></td>
</tr>
<tr>
<td>transverse</td>
<td>3.5</td>
</tr>
<tr>
<td>emittances</td>
<td>3.75 µm.rad</td>
</tr>
<tr>
<td>Longitudinal emittance (total)</td>
<td>1</td>
</tr>
<tr>
<td></td>
<td>2.5 eVs</td>
</tr>
</tbody>
</table>

Initial Design
The first challenge is the high beam brightness ($N_b/\varepsilon_x$), which is limited by the effect of the space charge induced tune spread $\Delta Q_x$ in the low energy machines. This has been addressed (i) in the PS-Booster (PSB), by dividing the intensity by a factor 2, using two PSB pulses to fill the PS, and (ii) in the PS, by increasing the injection energy from 1 GeV to 1.4 GeV kinetic.

Although these measures brought the expected results in the transverse phase space, early experiments [5] demonstrated the need for improving the longitudinal procedure, because:

- Longitudinal emittance blew up excessively during debunching-rebunching at 25 GeV in the PS,
- Beam losses were unavoidable at PS ejection (25 ns bunch spacing with 100 ns kicker rise-time),
- There was a risk of generating edge bunches with transverse tails,
- The “clean” generation of a bunch train shorter than the PS circumference was impossible.

Improved Longitudinal Procedure
Instead of using debunching-rebunching to change the time structure of the beam, a multiple splitting scheme is now employed.

Six PSB bunches (two PSB batches of 3 + 3 or 4 + 2 bunches) are captured on harmonic 7 in the PS. For the generation of the nominal LHC bunch train with 25 ns spacing in the PS, the bunches are then split in three at 1.4 GeV using appropriate amplitude and phase parameters on three groups of cavities operating on harmonics 7, 14 and 21 respectively [6]. Once captured on harmonic 21, the beam is accelerated up to 25 GeV where each bunch is split twice in two. During that process the RF harmonics number is changed to 42 (by an additional RF system at 20 MHz) and finally 84 (40 MHz). Each of the six original bunches has then been split in 12, and 72 bunches are created on harmonic 84 (40 MHz). The 80 MHz systems are finally pulsed before ejection to shorten the bunches to 4 ns so as to fit into the SPS 200 MHz buckets. The final bunch train contains 12 consecutive empty buckets, providing a gap of ~320 ns (13 x 25 ns) for the rise-time of the ejection kicker.

By making the additional RF system able to operate both at 20 and 13 MHz, and using only two splittings in two bunch trains with 75 ns spacing in the two bunch train can easily be generated. Moreover, shorter bunch trains can easily be prepared by injecting less than 6 PSB bunches.

However, this new scheme requires a more dense beam from the PSB: 72 LHC bunches are produced from 6 PSB bunches (instead of initially 84 from 8). Therefore the intensity per PSB bunch and consequently, due to the fixed emittance, the beam brightness have to be 14% higher than with the debunching-rebunching scheme. The ultimate beam is consequently more difficult for the PSB to provide.
Changes in the LHC machine itself with respect to the initial conceptual design [7] have a similar additional effect: the decision to increase the crossing angle of the beams in the interaction region has lead to another 10% increase of the intensity per bunch, and the change of the interaction point beta function $\beta^*$ from 0.5 to 0.55m has added another 5%.

**Experimental results**

Using the above-mentioned technique, beams with better than nominal characteristics are regularly obtained in the PS [2], and quasi-nominal performance is attained at the extraction energy of the SPS [1]. However, $\pm 10\%$ beam losses are unexpectedly observed between PS ejection and SPS at 450 GeV. Suspicion lies on the impedance of the SPS ejection kickers. Since more kickers have to be installed during the next years and no short term upgrade is possible, it is cautious to assume that the problem will remain at this level for a number of years.

**Performance Summary**

Table 2 summarizes the enhanced demand on the intensity per bunch from the PSB, within constant transverse emittances, which results from the above-mentioned effects. Transmission efficiencies from the PSB to the SPS at 450 GeV have been taken as 85% for nominal and 80% for ultimate beam.

Table 2: Protons per bunch from PSB (25 ns bunch train)

<table>
<thead>
<tr>
<th></th>
<th></th>
<th></th>
<th></th>
</tr>
</thead>
<tbody>
<tr>
<td>LHC nominal bunch</td>
<td>$1.00 \times 10^{11}$</td>
<td>$1.15 \times 10^{11}$</td>
<td>1.15</td>
</tr>
<tr>
<td>PSB nominal bunch</td>
<td>$10.5 \times 10^{10}$</td>
<td>$16.3 \times 10^{10}$</td>
<td>1.54</td>
</tr>
<tr>
<td>LHC ultimate bunch</td>
<td>$1.7 \times 10^{11}$</td>
<td>$1.00 \times 10^{11}$</td>
<td>1.00</td>
</tr>
<tr>
<td>PSB ultimate bunch</td>
<td>$17.9 \times 10^{10}$</td>
<td>$25.4 \times 10^{10}$</td>
<td>1.42</td>
</tr>
</tbody>
</table>

Since experiments have shown that the PSB can only deliver a maximum of $\approx 20 \times 10^{11}$ ppb within the transverse emittances required for the LHC, it is clear that:

- The nominal beam can be produced with a reasonable safety margin, as well as all other types of beams which are less demanding (75 ns bunch trains, single bunches, etc.),

- The ultimate beam cannot be obtained by the present means.

**NEEDS OF THE LHC UPGRADES**

Considering the long time needed for the R&D, it is necessary to start already today analysing the possible means to increase the LHC luminosity beyond the present ultimate level. After an initial “brain-storming” [8], this activity is now one of the main subject of the “H3” (High energy - High intensity - Hadron beams) network supported by the European Union inside the “CARE” (Coordinated Accelerator Research in Europe) integrated activity [9].

For short bunches the luminosity $L$ can be expressed as:

$$L = \gamma \frac{N_b^2 f_{rep} \Delta Q_{bb}}{4\pi\varepsilon_n \beta^*} F$$

where $f_{rep}$ is the average bunch repetition frequency, $\gamma$ the relativistic gamma factor and $F$ a form factor:

$$F = 1/\sqrt{1 + \left(\frac{\theta_c \sigma_z}{2\sigma^*}\right)^2}$$

$\theta_c$ is the crossing angle, $\sigma_z$ the r.m.s. bunch length and $\sigma^*$ the rms transverse beam size at the IP.

It is worthwhile remarking that the total linear tune shift $\Delta Q_{bb}$ for short bunches colliding with a crossing angle in alternating horizontal-vertical planes is also proportional to $F$:

$$\Delta Q_{bb} \propto \frac{N_b}{2\pi\varepsilon_n} F$$

**Schemes**

Many of the upgrades are based on increasing the crossing angle $\theta_c$ and sometimes $\sigma_z$ to make $F$ smaller. For a given $\Delta Q_{bb}$, brightness $N_b/\varepsilon_n$ and $N_b$ itself can then be increased, which results in a higher luminosity.

Another possibility is to increase the number of bunches, using a smaller bunch spacing. These being generated from an unchanged number of PSB bunches, the PSB brightness has to be directly proportional to the number of bunches.

Finally a more “exotic” solution is to use very long and “flat” bunches which can potentially provide a luminosity 1.4 times larger than Gaussian bunches. A sophisticated procedure has recently been proposed to generate such bunches [10] and it can be used to estimate the necessary brightness from the injectors.

**Beam Characteristics required from the injectors**

Using the terminology in reference [8], the proposals are divided in 3 phases of increasing cost and performance: phase 0 contains no hardware changes, phase 1 assumes hardware upgrades in the LHC insertions and/or in the injector complex, and phase 2 involves the replacement of major pieces of equipment.

The main scenarios considered are listed in Table 3. The required brightness with respect to the ultimate beam at the exit of the PSB is indicated in the fourth column. The bunch characteristics at ejection from the PS are given in the fifth column, assuming 85% transmission between PS at 25 GeV and LHC at 7 TeV.

It is clear from Table 3 that the ultimate beam parameters are an absolute minimum and that most LHC upgrade scenarios are even more demanding in terms of brightness from the injectors.
Table 3: Needs of the LHC luminosity upgrade scenarios (*assuming a transmission efficiency of 85 % between PS and LHC)

<table>
<thead>
<tr>
<th>Phase</th>
<th>L (10^{34} cm^{-2}s^{-1})</th>
<th>Comment</th>
<th>PSB brightness factor</th>
<th>N_p with 25 ns spacing from the PS*</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>3.6</td>
<td>Large crossing angle</td>
<td>1.5</td>
<td>3 \times 10^{11} (\varepsilon_n = 3 \mu m.rad)</td>
</tr>
<tr>
<td>1a</td>
<td>4.6</td>
<td>Low \beta^* + large</td>
<td>1.0</td>
<td>2 \times 10^{11} (\varepsilon_n = 3 \mu m.rad)</td>
</tr>
<tr>
<td></td>
<td></td>
<td>crossing angle +</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>high h RF</td>
<td></td>
<td></td>
</tr>
<tr>
<td>1b</td>
<td>7.7</td>
<td>Phase 1a + 15 ns bunch</td>
<td>1.7</td>
<td>3.4 \times 10^{11} (\varepsilon_n = 3 \mu m.rad)</td>
</tr>
<tr>
<td>1c</td>
<td>10</td>
<td>~ 80 long bunches @1.6 \Delta \nu</td>
<td>1.9</td>
<td>3.8 \times 10^{11} (\varepsilon_n = 3 \mu m.rad)</td>
</tr>
<tr>
<td>2</td>
<td>10</td>
<td>1 TeV injector + beam-beam compensation</td>
<td>&gt; 1</td>
<td>&gt; 4 \times 10^{11} (\varepsilon_n = 6 \mu m.rad)</td>
</tr>
</tbody>
</table>

POSSIBLE IMPROVEMENTS

Space charge induced tune-spread at low energy limits the brightness that the PSB and PS can deliver:

\[ \Delta Q_{sc} \approx \frac{N_p}{\varepsilon_n} \frac{1}{\beta^* \gamma^2} \]

In a given accelerator, the problem is most severe at low energy. The obvious possibility for improvements is therefore to increase the injection energy.

Another solution is to increase the number of protons per bunch at a higher energy, using longitudinal manipulations. Brightness is therefore increased in the same proportions, without suffering from an excessive \( \Delta Q_{sc} \). However, this results in a decrease of the number of bunches per machine pulse.

For the time-being the limitations of the SPS beyond the ultimate brightness have not been analysed.

**Batch Compression in the PS**

The following longitudinal beam gymnastics can be applied in the PS:

- Inject 7 (4 + 3) bunches from two PSB batches into adjacent h=9 buckets of the PS,
- Accelerate this beam up to an intermediate energy where space charge is reduced by more than 14/9,
- Compress the 7 bunches into 7/14 of the PS circumference by adiabatically increasing from h=9 to 10, 11, 12, 13, and 14,
- Resume acceleration on h=14 up to 25 GeV,
- Triple split the bunches using RF on h=28 and 42 (similar process as used at 1.4 GeV for the 25 ns bunch train [6]),
- Double split the bunches, changing the harmonic number from 42 to 84,
- Proceed with bunch compression before ejection using the combined action of RF on h=84 and h=168. This results in 42 bunches, populated with up to 2.6\times10^{11} protons and spaced by 25 ns, that can potentially be delivered to the SPS every 3.6 s.

The evolution of the bucket heights during the batch compression (Fig. 1) shows that the stable regions housing the 7 bunches stay large enough and that the process is smooth and can be made quasi-adiabatic.

Figure 1: Batch compression \( h=9 \rightarrow h=14 \) in the PS.

The filling scheme of the LHC has to be reviewed accordingly. More accelerator cycles are needed, and more kicker rise and fall times have to be accommodated, resulting in a slightly reduced number of bunches circulating in the collider. Among the many possible schemes, one interesting variant is shown in Figure 2. The main consequences are that the LHC is populated with 7\% less bunches (2604 instead of 2808), and the filling time is increased by 33 \% (the duration of the SPS cycle goes from 21.6 s to 28.8 s).

Figure 2: LHC filling with batch compression in the PS.
New Injector for the PSB (Linac4)

The energy of the beam injected in the PSB can be increased to 160 MeV (which corresponds to doubling $\beta^p$) by replacing the present Linac2 with a new accelerator called Linac4. This new linac, which is capable of becoming the front-end of a high energy / high power Superconducting Proton Linac (SPL), is presently under study. Components are being developed with the support of the European Union inside European laboratories [9] and of the International Science and Technology Center (Moscow) in Russia [11].

The basic layout of this ~90 m long machine is shown in Figure 3.

![Figure 3: Schematic layout of linac4](image)

Beam dynamics studies and computer simulations have indeed shown that the brightness of the PSB beam can be expected to double [12].

The following scheme will be used in the longitudinal phase plane:

- The PSB captures and accelerates beam on $h=3$.
- Before transfer to the PS, the distance between bunches is reduced by applying $h=1$.
- 12 (4×3) bunches are injected from a single PSB batch into adjacent $h=14$ buckets in the PS.
- Acceleration takes place on $h=14$ up to 25 GeV.
- Bunches are triple split using RF on $h=14$, 28 and 42 (similar process than used at 1.4 GeV for the 25 ns bunch train [6]).
- Bunches are again split in two, changing the harmonic number from 42 to 84.
- Bunch compression is effected before ejection using the combined action of RF on $h=84$ and $h=168$.

This way, 72 bunches populated with up to $2\times10^{11}$ protons and spaced by 25 ns can potentially be delivered to the SPS every 2.4 s.

Obviously a batch compression process could be applied additionally, which would make the PS able to deliver 48 bunches populated with up to $3\times10^{11}$ protons every 2.4 s.

SPL replacing the PSB

In its present design, the SPL [13] provides an H’ beam at a kinetic energy of 2.2 GeV which is ideally suited for accumulation in the PS, replacing the PSB as injector. For a $\Delta Q_{\text{acc}}$ at injection of 0.31, 72 bunches of $4\times10^{11}$ protons with 25 ns spacing could be delivered by the PS.

The operation of the PS will become very simple and flexible, because:

- The longitudinal distribution of the beam can easily be tailored to the needs (number of bunches, distance between bunches, …).
- No RF gymnastics is needed.
- Dead-time between PS cycles can be minimized, by virtue of the high repetition rate of the SPL (50 Hz), so that the cycling rate of the PS is optimized.

Between 1 and 80 bunches populated with up to $4\times10^{11}$ protons and spaced by 25 ns can potentially be delivered to the SPS every 1.8 s.

Rapid Cycling Synchrotron replacing the PSB

In case the high beam power capability of the SPL is not needed, it would be worthwhile considering the replacement of the PSB by a higher performance synchrotron. A typical goal could be to reach in the PS twice the ultimate brightness, like with the SPL. This determines the maximum energy, which has to be at 2.2 GeV.

Space charge at injection in this new synchrotron should be compatible with the brightness that the PS accepts at 2.2 GeV. Assuming it has the same size as the present PSB, one can derive that the injection energy has to be approximately 360 MeV.

To minimize the duration of the PS injection flat porch, the new synchrotron must either have 4 rings or be fast cycling (typically 50 Hz).

To match the SPL potential it must be able to deliver $8.4\times10^{12}$ protons per pulse within transverse emittances of $\varepsilon_t=2.5\ \mu\text{m}.\text{rad}$.

Under these conditions, the same beam as attained with the SPL could be delivered to the SPS, namely up to 80 bunches populated with up to $4\times10^{11}$ protons and spaced by 25 ns.

Comparative summary

The solutions envisaged for the LHC luminosity upgrade are summarised in Table 4. They differ enormously in cost and in performance potential.

The “batch compression” in the PS is “low cost” and can be quickly implemented. It is the ideal method for exploring the potential and limitations of the SPS and LHC beyond the ultimate brightness. However, (i) its sophistication makes it difficult to maintain and a risk for reliability, (ii) it lengthens the LHC filling time and (iii) the luminosity is reduced by approximately 10% because of the smaller number of bunches in LHC.

Linac4 is a “middle cost” solution to reliably deliver the beam, and reach the ultimate parameters inside the LHC in regular operation. Combined with batch compression in the PS it will allow to explore further the capabilities of SPS and LHC.

The replacement of the PSB is the most expensive possibility. The choice between SPL and RCS will be determined by arguments external to the LHC upgrade. These solutions are attractive for their high level of beam performance, their high potential and flexibility for the next decades and the improved reliability resulting from the replacement of ~40 years old accelerators.
ACKNOWLEDGEMENTS

The results and the ideas presented in this paper have evolved during the last 10 years, and it is impossible to do justice to all the colleagues who have contributed. Nevertheless, special credit is due to K. Schindl, J. Gareyte and F. Ruggiero for the efficient brain-storming spirit they have managed to create.

REFERENCES


Table 4: Comparative summary of the possible changes in
the injectors

<table>
<thead>
<tr>
<th></th>
<th>Batch compression in the PS</th>
<th>Linac4</th>
<th>RCS</th>
<th>SPL</th>
</tr>
</thead>
<tbody>
<tr>
<td>Cost (P+M) [MCHF]</td>
<td>&lt; 1</td>
<td>~ 70</td>
<td>~ 200</td>
<td>~ 500</td>
</tr>
<tr>
<td>Delay [years]</td>
<td>~ 2</td>
<td>3-4</td>
<td>4</td>
<td>5-6</td>
</tr>
<tr>
<td>Operation: Comfort</td>
<td>Delicate</td>
<td>+</td>
<td>++</td>
<td>+++</td>
</tr>
<tr>
<td>Reliability</td>
<td>Limited</td>
<td>+</td>
<td>++</td>
<td>+++</td>
</tr>
<tr>
<td>Nb. of bunches per PS pulse</td>
<td>42</td>
<td>72 (48)</td>
<td>72</td>
<td>1-80</td>
</tr>
<tr>
<td>Intensity per PS bunch [p/b]</td>
<td>2.6×10^{11}</td>
<td>2×10^{11} (3×10^{11})</td>
<td>4×10^{11}</td>
<td>4×10^{11}</td>
</tr>
<tr>
<td>Repetition period [s]</td>
<td>3.6</td>
<td>2.4</td>
<td>1.8</td>
<td>1.8</td>
</tr>
<tr>
<td>BEST USE</td>
<td>Exploratory tests</td>
<td>Reliable operation + 40% of upgrades</td>
<td>Reliable operation + all upgrades</td>
<td>Reliable operation + all upgrades</td>
</tr>
</tbody>
</table>

Acknowledgements

The results and the ideas presented in this paper have evolved during the last 10 years, and it is impossible to do justice to all the colleagues who have contributed. Nevertheless, special credit is due to K. Schindl, J. Gareyte and F. Ruggiero for the efficient brainstorming spirit they have managed to create.
Abstract
To prepare the SPS as injector to the LHC the SPS impedance was significantly reduced. The results of this programme followed by beam measurements are presented together with present intensity limitations.

INTRODUCTION
Recently an LHC beam with nominal intensity and bunch parameters was accelerated in the SPS to top energy, 450 GeV [1]. This is the result of an SPS upgrade [2] followed by a few years of extensive machine development (MD) studies.

The SPS impedance reduction programme was an important element in the SPS upgrade together with other hardware modifications.

Measurements with the beam done following completion of the impedance reduction programme in 2001 have shown significantly improved single bunch stability and absence of the uncontrolled emittance blow-up due to microwave instability at least up to ultimate intensity [3]. However to obtain the nominal LHC beam at top energy many other problems, mainly connected with e-cloud, beam losses, beam loading and coupled bunch instabilities, had also to be solved.

SPS IMPEDANCE
Impedance reduction
For many years the longitudinal single bunch instability was one of the most serious intensity limitations. The main sources of impedance causing this instability were found from measurements of the unstable spectrum of long single bunches injected into machine with RF off [4]. In preparing the old machine in its new role as injector of LHC this impedance was significantly reduced.

The pumping ports (around 1000) were shown to be the main source of the microwave instability (responsible for the signals above 1.3 GHz). Depending on the frequency their total R/Q varies in the range (25 - 45) kOhm, Q~ 50. Now 98% of them are shielded. During the 5 months of the 2000/2001 shutdown 400 main dipole magnets were displaced to install 2400 sub-assemblies of 30 different screen models [5].

The 400 MHz instability sources were not initially clear. 16 septa MSE and MST were shielded during 1998-2000 (B. Goddard, A. Rizzo), but the instability was still there. During the 2000/2001 shutdown 6 MKE and MKP kickers were removed or screened and the strong signal at 400 MHz disappeared. Five MKE kickers with screens were again installed in the ring in 2003, four more kickers will be installed in 2006.

Removal of lepton cavities (3 RF systems) and other non-used equipment helped as well in the total impedance reduction.


Known SPS impedance
The present impedance budget includes:
- two TW RF systems
  - main 200 MHz RF system consisting of two cavities of four sections and two cavities of five sections,
  - the fourth harmonic RF system at 800 MHz consisting of two cavities, used for Landau damping;
- five MKE and three MKP kickers with RF by-passes;
- resistive-wall impedance of stainless steel chamber;
- e-cloud.

Impedance measurements with the beam
In the longitudinal plane the following measurements were performed to refine the SPS impedance model.
- The spectrum of long unstable bunches, used in the past to identify resonant impedances with high R_{sh}/Q and low quality factor Q [4], is used now to monitor the changes of impedance.
- Measurements of bunch lengthening with intensity provide information about ImZ/n below the microwave instability threshold and about high frequency impedances for intensities above this threshold [3].
- Measurements of the coherent frequency shift with intensity allow the changes of ImZ/n with time to be seen [3, 7].
- The stable phase shift measured as a function of intensity allowed the value of ReZ to be estimated [8].
- Measurements of the unstable mode spectrum of the LHC beam help to identify HOMs with high R_{sh} and quality factor Q responsible for coupled-bunch instabilities observed at high energies during the cycle.
Both imaginary and real parts of the effective impedance have a strong dependence on bunch length $\tau$. Measurements of synchrotron tune shift were performed below and above transition, at 14 GeV/c and 26 GeV/c. Reasonable agreement between calculations and measured quadrupole frequency shift as a function of bunch length and intensity at 14 GeV/c was found using effective impedances $\text{Im}Z_{\text{eff}}(\tau)$ and $\text{Im}Z_{\text{eff}}^2(\tau)$ [9]. The contribution to the measured frequency shift from the coherent part is significantly smaller than from the incoherent part, so that their ratio is less than 0.15. The imaginary part of the SPS impedance up to 1 GHz used in calculations of the quadrupole frequency shift is shown in Fig. 1. It corresponds to the low frequency inductive impedance $\text{Im}Z/n \approx 7.5$ Ohm.

Measurements of synchronous phase shift with intensity allowed the dependence of energy loss $U$ on bunch length $\tau = 4\sigma$ to be determined experimentally [8]. The bunch length variation was provided by changes of capture voltage in the SPS in the range (0.6-3.0) MV. For a Gaussian bunch the energy loss per turn and per particle is

$$U_n = e^2 N \frac{\omega_0}{\pi} \sum_{\mu=0}^{\infty} \text{Re}Z_n(p\omega_0)e^{-(p\omega_0\sigma)^2},$$

where $N$ is bunch intensity and $\omega_0$ is the circular revolution frequency. Main contributions come from four cavities of the 200 MHz RF system and five MKE kickers. The real part of these impedances is presented in Fig. 2. Contribution from resistive wall impedance for $N \approx 10^{10}$ is $\bar{U} = 0.8$ KeV for $\sigma = 0.6$ ns. Measurements and calculations based on this impedance model are in very good agreement [8].

**In the transverse plane** a series of beam measurements were performed as well to monitor the SPS impedance change and to refine the impedance model. This includes measurements of

- coherent tune shifts with intensity ($\text{Im}Z_T$) [6],
- growth rates as a function of chromaticity ($\text{Re}Z_T$) [6],
- the fast transverse instability threshold ($Z_V$) [11],
- betatron phase beating to determine transverse impedance localisation [12].

An impedance model which gives a reasonable fit to the growth rate dependence on chromaticity and coherent tune shifts with intensity is based on a broad-band resonator with resonant frequency $f_r = 1.3$ GHz, $Q=1$ and $Z_V \approx 10$ MOhm/m in the vertical plane and $Z_H \approx 7$ MOhm/m in the horizontal plane. Measurements of coherent tune shifts with intensity indicated a 40% impedance reduction in 2001 and a 50% increase in 2003, after re-installation of five MKE kickers in the ring [13]. The vertical impedance of one MKE kicker can be approximated by a broad-band resonator centered at 2 GHz with peak value 0.6 MOhm/m [14]. The contribution of five kickers to the imaginary part of the vertical impedance is then $\approx 3$ MOhm/m.

**INTENSITY LIMITATIONS**

**Single bunch**

The nominal bunch intensity at the top energy in the LHC is $1.15 \times 10^{11}$ and the ultimate intensity is $1.7 \times 10^{11}$. 
To allow $\sim 10\%$ particle losses in the SPS and LHC
$1.3 \times 10^{11}$ should be injected into the SPS to achieve
nominal value in the LHC. Nominal injected emittance is
0.35 eVs and $4\sigma$ bunch length is 4.2 ns.

On the 26 GeV/c flat bottom the quadrupole oscillations
of bunches with longitudinal emittance of $\varepsilon = 0.2$ eVs in-
jected into a mismatched voltage are not damped above an
intensity $3 \times 10^{10}$. Some slow (growth rate $\sim 2$ s) lon-
gitudinal instability is observed for intensities more than
$4 \times 10^{10}$.

Fast transverse instability of the single proton bunch was
observed in 2002 [15] and studied in more detail in 2003
[11]. Bunches with $\varepsilon = 0.3$ eVs, $\tau = 3.8$ ns and intensity
above $1.2 \times 10^{11}$ had strong losses at injection for zero chromaticity. Increase of chromaticity (also used in
operation to prevent a vertical instability due to e-cloud)
cures this instability. Nevertheless the threshold, which in-
creases with bunch length and is probably now for the LHC
bunch just below ultimate intensity, could be close to nom-
inal intensity in 2006, when more MKE kickers will be in-
stalled. These estimations are obtained by calculations with
MOSES code (Y. H. Chin, 1988). More accurate predic-
tions will be obtained using the Head Tail code [16] which
can take into account the flat geometry of the SPS vacuum
chamber and the space charge effect.

The LHC beam

The LHC beam in the SPS consists of 3-4 batches, each
of 72 bunches, injected at 2.4 s intervals. Bunches are
spaced at 25 ns, while the RF period is 5 ns.

To prepare operation with this very high intensity beam
the 200 MHz RF system was upgraded [17]:

- one feedback system per cavity,
- new feed-forward system,
- new longitudinal damper with (0 - 3 MHz) bandwidth,
- new 1 MW couplers.

The bandwidth of transverse damper was increased to
20 MHz and the gain was also upgraded [18].

Present intensity limitations for the LHC beam could be
engineered as follows.

- e-cloud.
- Injection: capture loss.
- Acceleration: coupled bunch instabilities.
- Flat top: requirements for injection to LHC.

Intensity effects and limitations due to e-cloud in the SPS
deserve special attention and have been discussed in other
talks at this Workshop. Below is a short summary of limi-
tations and cures [1].

Figure 3: Dependence of relative beam loss on batch intensity measured in 2002 (group of three circles), 2003 with
1 (black circle), 2 (red), 3 (blue) and 4 (green) batches in
the ring and in 2004 (old and new working points). 2 MV
voltage at 26 GeV.

- e-cloud leads to transverse emittance blow-up and insta-
abilities:
  - coupled bunch in H-plane (a few MHz),
  - single bunch in V-plane affecting the tail of the batch
    ($\sim 700$ MHz).
- Scrubbing run increases the threshold from $0.3 \times 10^{11}$
to $1 \times 10^{11}$.
- Transverse feedback helps to damp coupled-bunch
  modes in H-plane with growth rates up to 40 turns.
- High chromaticity is used as a cure for V-plane.

Injection

In 2002 an LHC beam with nominal intensity and lon-
gitudinal parameters was accelerated in the SPS to top
energy, 450 GeV. In 2003 even higher intensities were
achieved. However relative capture losses increased by al-
most 30% for the same batch intensity and capture voltage
(Fig. 3) [19]. To obtain a nominal intensity LHC beam,
$1.15 \times 10^{11}$ p/bunch at 450 GeV, 15% more particles had
to be injected into the SPS when using a 2 MV capture volt-
age. For smaller voltages capture losses were even higher
reaching 40% for the match voltage of 750 kV (Fig. 3). In
2003 and 2004 a significant amount of machine develop-
ment (MD) study time in the SPS was devoted to this beam
loss problem.

Losses were measured at the end of the 10.86 s long
injection plateau (26 GeV) and after the start of acceleration
(30 GeV) with the beam current transformer (BCT). As can
be seen from Fig. 3 capture losses depend strongly on the
batch intensity $N$ so that the relative loss is proportional to
$N$ and the absolute loss to $N^2$. In these measurements with
a voltage of 2 MV the number of batches did not have any
evident influence on the beam losses.

Relative losses were significantly reduced for 75 ns
bunch spacing and nominal intensity per bunch (5% loss
for 16 bunches with $1.2 \times 10^{11}$/bunch).
It is also seen that injection losses have an asymmetric character; practically all lost particles moved away from the front of the batch rather than from the back. Above transition energy this implies that lost particles have or acquire some negative energy deviation.

The asymmetric character of the beam loss can be explained by particles having energy loss due to the resistive impedance in the machine [8] and even possibly due to electron-cloud production. An energy loss \( U \) per particle leads to an accelerating-type bucket with a synchronous phase \( \phi_s \simeq U/(eV) \) on the flat bottom. As a result gaps exist between the buckets and particle trajectories outside the bucket lead eventually to lower energy. The azimuthal size of the gap between buckets is \( \delta \phi \simeq 2\sqrt{\pi} \sin \phi_s \). For a 0.5 ns gap one needs \( \phi_s = 1.8 \) deg at 200 MHz and \( U/(eV) \simeq 0.03 \) or only \( U = 60 \) keV for \( V = 2 \) MV.

Increase of voltage at injection helps to reduce losses (2 MV voltage is used instead of matched voltage of 750 kV). Losses are even less when 2 MV is increased to 3 MV shortly (within 100 ms) after injection, see Table 1. This scheme gives loss reduction for one batch, but could produce satellite bunches when voltage dips are used for subsequent batches due to loss and recapture of particles.

Bunches injected from PS have bunch length variation in the range \((4.2 \pm 0.5) \) ns. Losses are reduced for shorter bunches with smaller emittance, but then the beam is more unstable on the flat bottom.

<table>
<thead>
<tr>
<th>( V_{inj} ) (MV)</th>
<th>( V_{fb} ) (MV)</th>
<th>( N_{mean} ) ( \times 10^{12} )</th>
<th>SD ( \times 10^{12} )</th>
<th>loss ( % )</th>
<th>SD ( % )</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.0</td>
<td>2.0</td>
<td>33.23</td>
<td>0.26</td>
<td>9.1</td>
<td>0.9</td>
</tr>
<tr>
<td>3.0</td>
<td>3.0</td>
<td>33.18</td>
<td>0.29</td>
<td>9.5</td>
<td>1.0</td>
</tr>
<tr>
<td>2.0</td>
<td>3.0*</td>
<td>33.18</td>
<td>0.43</td>
<td>8.4</td>
<td>1.1</td>
</tr>
<tr>
<td>2.0</td>
<td>3.0</td>
<td>8.55</td>
<td>0.11</td>
<td>7.8</td>
<td>0.6</td>
</tr>
<tr>
<td>2.0</td>
<td>3.0</td>
<td>8.67</td>
<td>0.09</td>
<td>4.4</td>
<td>0.5</td>
</tr>
</tbody>
</table>

Table 1: Beam losses measured in 2004 for old (above double line) and new (below double line) working points for different voltages at injection \( V_{inj} \) and on the flat bottom \( V_{fb} \). Voltage programme during the ramp corresponds to 0.6 eVs bucket area. The first three lines - 4 batches in the ring, the rest - one batch in the ring. Comments: (*) - 4 dips to 2 MV on the flat bottom. Nominal vertical chromaticity of 0.365. The 800 MHz RF system is used on the flat bottom (with 235 kV) and during the cycle (600 kV) in bunch shortening mode.

Apart of capture loss, a continuous flux of particles from the bucket along the flat bottom, after injection, was observed [19]. In 2004 losses along the flat bottom were reduced when the working point was changed from (26.19, 26.13) to (26.13, 26.19) [20].

**Acceleration**

Single batch with \( 2 \times 10^{10}/\)bunch is unstable at \( \sim 280 \) GeV (16 s). This threshold was \( 3 \times 10^{10}/\)bunch in 2002 and most probably has decreased due to the re-installation of MKE kickers in the ring which causes increase of the inductive part of the SPS impedance and corresponding loss of Landau damping.

The 800 MHz RF system is used in bunch-shortening (bs) mode during the whole cycle to increase synchrotron frequency spread (almost by factor 5) and stabilise the beam. This is still not sufficient for stability of the nominal beam on the flat top, so preventive emittance blow-up is applied. Beam excitation on the ramp (at 15 s) with band-limited noise around \( 2f_s \) on the 200 MHz voltage amplitude is used in addition to the emittance blow-up from the filamentation in the mismatched voltage at injection (2 MV instead of 700 kV gives \( \sim 0.4 \) eVs).

Measurements of unstable beam spectrum during the development of this instability suggest that the HOMs in the 200 MHz RF system at 629 MHz and 912 MHz could be responsible.

**Extraction**

At extraction in the SPS the bunch length should be significantly below 2.5 ns and the emittance less than 0.7 eVs to minimise losses at injection into LHC where the 400 MHz RF system is used for beam capture. Four batches of 72 bunches with \( 4\sigma \) average bunch length of \( 1.6 \pm 0.2 \) ns were obtained at 450 GeV for an intensity of \( 1.15 \times 10^{11}/\)bunch with voltage \( V_{200} = 7 \) MV and \( V_{800} = 0.7 \) kV which corresponds to \( \varepsilon_{2\sigma} = 0.6 \) eVs. With a stable beam the bunch to bunch phase on the flat top is inside 130 ps, see Fig. 4, [21].

Note that \( 2\sigma \) emittance contains only \( \sim 85\% \) of particles for a Gaussian bunch limited at separatrix. For nom-
inal intensities 6σ bunch length (2.4 ns) is at the limit of the bucket size in the LHC (2.5 ns). It is possible that the 200 MHz capture system in the LHC will be required for ultimate intensities especially if the SPS impedance will be further increased.

**SUMMARY**

The main intensity limitations in the SPS for the LHC beam are:

- Intensity dependent capture loss (~ 8%).
- Although coupled bunch instabilities are cured,
  - the 800 MHz in bs-mode increases the peak line density (problem for e-cloud, MKE heating).
  - emittance blow-up could lead to extra losses at injection into LHC with 400 MHz bucket.
- Beam loading in the 200 MHz and 800 MHz (efficiency as a Landau cavity) RF systems.
- MKE heating due to its resistive impedance [22].
- e-cloud and possibly fast transverse instability with more MKE kickers or higher bunch intensities. A cure using chromaticity in conjunction with high voltage increases losses.

Possible improvements include:

- Further SPS impedance reduction (MKE shielding [23], improved passive damping of HOMs at 629 MHz and 912 MHz, search for transverse impedances...)
- Shorter bunches from PS with the same or larger emittance (requires extra RF voltage in the PS).
- Increased voltage from the 800 MHz RF system (one more cavity in operation in 2005).
- Emittance blow-up to increase the threshold of coupled bunch instability on the flat top $\propto \varepsilon^2$ (0.75 eVs needed for ultimate intensity).
- A 200 MHz RF system in the LHC for capture.
- Capture loss studies (RF noise, e-cloud, machine resonances...)

**REFERENCES**

Multiturn Extraction Based on Trapping in Stable Islands


Abstract

For some applications an intermediate extraction mode between fast (one turn) and slow (several thousand turns) is needed. This is the case of the five-turn extraction used to transfer the proton beam from the CERN Proton Synchrotron (PS) to the Super Proton Synchrotron (SPS). Unavoidable losses and poor betatron matching with the receiving machine affect the present approach, which is based on beam slicing by means of an electrostatic septum. These features are rather serious obstacles to an intensity upgrade of the PS/SPS Complex. To overcome these difficulties, a novel extraction technique was proposed recently. By using nonlinear magnetic elements, stable islands can be generated in the transverse phase space. Furthermore, provided the linear tune is varied slowly, it is possible to trap the charged particles inside the stable islands in order to split the beam into different beamlets. Once generated, the distance between these beamlets can be tailored to the extraction needs by simply increasing the distance of the tune from the resonance value. The results of numerical simulations for the novel technique are presented and discussed in detail together with the outcome of the intense experimental studies performed to assess the validity of this approach.

INTRODUCTION

Since the approval of the CERN Neutrino to Gran Sasso Project (CNGS) [1] and the subsequent efforts devoted to a feasibility study of an intensity upgrade of the PS/SPS complex [2], the special extraction mode, the so-called Continuous Transfer (CT) [3], was reviewed. Such an extraction scheme is required to minimise the filling time of the SPS at 14 GeV/c, while reducing the beam emittance so to overcome the aperture limitation at SPS injection in the vertical plane\(^1\). The CT extraction was developed in the seventies [3] with the aim of extracting the beam from the PS in five consecutive turns using an electrostatic septum to slice the beam in the horizontal plane, the tune being 6.25 (see Fig. 1 for a sketch of the principle). The main drawbacks of this technique are the intrinsic losses on the electrostatic septum and the poor betatron matching of the five slices, which might transfer into injection losses in the SPS [5].

Recently, another alternative method was proposed, where the beam is split in the transverse phase space by means of adiabatic capture inside stable islands of the fourth-order resonance [6]. The method was then generalised by using other stable resonances [7]. On the experimental side, intense efforts were devoted to the demonstration of such a novel technique since the year 2002, when beam splitting was observed using a low-intensity single-bunch beam [8]. The key issue, i.e. whether the method would work for a high-intensity bunch, was considered during the years 2003 and 2004 (see Refs. [9, 10] for more details).

THE NOVEL MULTITURN EXTRACTION

In the novel approach [6] nonlinear elements such as sextupoles and octupoles are used to generate stable islands in transverse phase space. By varying the horizontal tune, particles can be selectively trapped in the islands by adiabatic capture: some will remain in the phase space area around the origin, while others will migrate to the stable islands. As a result, the beam is split into a number of parts in transverse phase space determined by the order of the resonance used, without any mechanical device. The separation between the islands can be controlled so that enough space is available for the beam to jump over a septum blade with almost no particles lost. A simple model representing the horizontal betatron motion in a circular machine under the influence of sextupole and octupole magnets was used (the motion in the vertical plane can be safely neglected). By assuming that the nonlinear magnets are located at the same place and the single-
kick approximation [11] is used, the one-turn transfer map can be expressed as \( x_{n+1} = M_n(x_n) \):
\[
\begin{pmatrix}
    x_{n+1} \\
    x'_{n+1}
\end{pmatrix} = R(2\pi \nu_n) \begin{pmatrix}
    x_n \\
    x^2_n + \kappa x^3_n
\end{pmatrix},
\]
where \((x, x')\) are obtained from the Courant-Snyder coordinates \((X, X')\) by means of the non-symplectic transformations [11]
\[
(x, x') = \frac{K_2 \beta H^{3/2}}{2} (X, X') \quad K_m = \frac{L}{B_0 \rho} \frac{\partial m B_0}{\partial x},
\]
where \(K_{2,3}\) the integrated sextupole (octupole) gradient, \(\beta H\) the value of the horizontal beta-function at the location of the nonlinear elements. \(R(2\pi \nu_n)\) is a \(2 \times 2\) rotation matrix of angle \(\nu_n\), the fractional parts of \(Q_H\). The map \(M_n\) is a time-dependent system through the linear tune. The time-dependence allows varying the phase space topology, thus creating and moving the islands. Also, it allows trapping particles inside the islands. A slow variation of the linear tune, adiabatic with respect to the time scale introduced by the betatron oscillations, pushes particles to cross the separatrix and to be trapped inside the newly-created islands. The trapping process has been simulated by using the model \(M_n\) with \(\kappa = -1.6\), while the tune \(\nu_n\) is varied according to the curve shown in Fig. 2. In the first part, the linear tune is decreased linearly from its initial value of 0.252 to 0.249. During this part, the capture process takes place. Then, a second linear decrease to the value 0.245 is performed which allows the islands to be moved towards higher amplitudes before extraction.

A set of Gaussian-distributed initial conditions has been generated, and its evolution under the dynamics induced by the map \(M_n\) is shown in Fig. 3. The trapping process is clearly visible in the picture: it generates five beamlets, well-separated at the end of the process. No particle is lost during the trapping phase, nor when the islands are moved.

![Figure 2: Linear tune \(\nu\) as a function of \(n\) for the fourth-order resonance.](image)

Figure 3: Evolution of the beam distribution during the trapping process with four islands. The different plots correspond to tune values represented by solid squares in Fig. 2. Each plot represents \(4.9 \times 10^5\) points. The initial Gaussian distribution is centred on zero and has \(\sigma = 0.045\). The co-ordinate system is the one defined in Eq. (2).

Not only the five beamlets have similar surfaces, but also their shape matches the phase space topology very well, making the five parts similar as far as transverse properties are concerned. It is worthwhile pointing out that the first four extracted beamlets have exactly the same emittance, as their shape is dictated by the same phase space structure, i.e. the island along the positive \(x\) axis. In this respect, the novel approach proves to be superior to the present CT extraction mode.

**MEASUREMENT CAMPAIGN**

**Machine and Instrumentation**

In the PS machine dedicated sextupoles and octupoles have been installed in sections 21 and 55 (two sextupoles in each section) and 20, 21 (two octupoles in series in section 20 and one in section 21) to generate the stable islands. Their location is optimised to have at least one island with zero angle at the extraction septum and extraction kicker. Only one set of sextupoles and octupoles can be used at a time. During the measurement campaign in 2002 [8] sextupoles in 21 and octupoles in 20 were used. In 2003, mainly the sextupoles in section 55 and the octupoles in section 20 were powered, while in the year 2004 sextupoles in section 21 and the single octupole in the same section are used. It is worthwhile mentioning that the PS lattice features the minimum of \(\beta_H\) in even straight sections. A sketch of the PS circumference together with the
The key elements used in the experiments are shown in Fig. 4. Phase space reconstruction and beam profile measurements are the techniques applied in the experimental study. The first ensures that the right phase space topology is generated by the nonlinear magnetic elements, while the latter is meant to record the evolution of the beam distribution during the trapping process and subsequent beam manipulations. The phase space reconstruction is based on turn-by-turn acquisition of the beam trajectory on two pickups 90° apart [8]. The beam profile is measured by means of a wire scanner [12]. Among the four installed, two for each transverse plane, the horizontal one in section 54 is routinely used for the measurements reported here. A scintillator is used to detect the secondary particles generated by the beam-wire interaction, thus reconstructing the beam profile. A summary of beam parameters is reported in Table 1.

2003 Measurement Campaign

Phase Space Measurement The low-intensity pencil beam is used for phase space reconstruction to avoid as much as possible beam filamentation. The beam trajectory is perturbed by the kicker magnet normally used for fast beam extraction and betatron oscillations are observed on two pickups 90° apart.

<table>
<thead>
<tr>
<th>Comments</th>
<th>Int. (10^{10})</th>
<th>$\epsilon_H(\sigma)$ (µm)</th>
<th>$\epsilon_V(\sigma)$ (µm)</th>
<th>$\Delta p/p(\sigma)$ (10^{-3})</th>
</tr>
</thead>
<tbody>
<tr>
<td>pencil beam</td>
<td>40</td>
<td>1.7</td>
<td>1.55</td>
<td>0.25</td>
</tr>
<tr>
<td>low-intensity</td>
<td>45</td>
<td>9</td>
<td>2.38</td>
<td>0.25</td>
</tr>
<tr>
<td>high-intensity</td>
<td>600</td>
<td>13.2</td>
<td>7.6</td>
<td>0.6</td>
</tr>
<tr>
<td>high-intensity a</td>
<td>600 b</td>
<td>9.4</td>
<td>6.4</td>
<td>0.7</td>
</tr>
</tbody>
</table>

a 2003 measurement campaign.

The pickups selected are those in sections 63 and 67 as they are in phase with the extraction septum and the extraction kicker used to excite the beam for the phase space reconstruction. An example of phase space measurement with a clear signature of stable islands is shown in Fig. 5. Beam position oscillations right after the kick indicate that the beam is rotating around the island’s centre and filamentation occurs until such oscillations are completely damped. The slow variation over time of the islands’ position is very likely due to particles’ diffusion outside the islands induced by longitudinal motion. The first few turns can be used to compute the so-called secondary frequency $\omega_{sec}$ [11], i.e. the betatron frequency around the fixed point. An analytical estimate can be obtained by computing the zero-amplitude frequency for the pendulum-like Hamiltonian [13]. The strong component originating from the fourth-order resonance is removed from the measured frequency spectrum by applying a stroboscopic analysis. Then, a refined algorithm to compute the tune [14] is applied. The large set of available data series has been used to derive an average value of the secondary frequency, thus yielding $\omega_{sec} = (0.0375 \pm 0.0008)$, which is in excellent agreement with the theoretical value $\omega_{sec} = 0.039$. A refined analysis was also attempted by taking into account the information concerning the different oscillation amplitude inside the island of the various measurement sets. This allowed measuring a detuning curve inside the islands. These data have been compared with numerical simulations and the outcome is plotted in Fig. 6. The good agreement is clearly visible, in spite of the measurement errors (probably underestimated). The straight lines represent the linear fit to the data. In the case of the experimental points, a fit weighted with the estimated errors has been applied, yielding $\omega_{sec} = (0.0408 \pm 0.0009)$, which is even closer to the theoretical result.

Adiabatic Capture of Low-Intensity Bunch The first step in the proof of principle of the novel multiturn extraction is the adiabatic capture of a low-intensity beam. To this aim, a beam with a large horizontal emittance, so to simulate the high-intensity beam, and a small vertical emittance to avoid nonlinear coupling between the two transverse planes, is used. No particle loss occurs during the capture and transport process [8]. A series of measurements was performed to assess the actual reversibility of the process by crossing the resonance twice, so to split the beam and merge it back. The reso-
Adiabatic Capture of High-Intensity Bunch The most difficult part was the capture of a high-intensity bunch of similar characteristics as those required for the proposed intensity upgrade for the CNGS [2]. Indeed, adiabatic trapping conditions were successfully established even for a bunch of intensity up to $625 \times 10^{10}$ protons (the nominal intensity being $600 \times 10^{10}$ protons). An example of the beam profile at the end of the capture process is shown in Fig. 7 (right), where both cases, i.e. nominal and record intensity, are plotted. Contrary to the low-intensity case, about 20% of the total beam is lost at the end of the capture process. It was conjectured that this effect might be induced by the strong vertical perturbation generated by the octupoles located at a high-$\beta_V$ section, which is imposed by mechanical constraints.

Detailed analysis of these data can be found in Ref. [10].

2004 Measurement Campaign Based on the experience gained in previous measurement campaigns, an additional octupole was installed during the 2003/2004 shutdown and it was used together with the two sextupoles located in the same section. The activities of phase space reconstruction and capture with the low-intensity beam were successfully resumed, thanks to a very good reproducibility of the PS machine over the years. This also indicates that the proposed approach features a remarkable robustness. The beneficial effect of the new octupole location in a low-$\beta_V$ section was reflected in a strong reduction of the losses in the case of capture of the high-intensity bunch: the optimisation of the working point and a proper setting-up of the longitudinal parameters reduced the losses down to $5 \sim 10\%$. However, the real improvement, leading to losses below the sensitivity of the beam transformer, was the dynamical change of the octupole during the capture process as seen in Fig. 8, where the evolution of the current of the key elements is shown. Indeed, by changing the octupole the beamlets can be separated while reducing the interval of tune variation. In Fig. 9 the evolution of the horizontal beam profile at the end of the capture process is reported. The fraction of particles in the beamlets is typically about $13 \sim 14\%$.

Single-turn extraction tests were also performed to assess whether the beamlets could be transported in a transfer line. To reduce losses at the extraction septum the beam-
Jets’ separation was reduced and the beam spot recorded in the transfer line via an Optical Transition Radiation (OTR) monitor (see Fig. 10). In the $(x, y)$ space two beamlets are projected onto the beam core, while the two lateral beamlets are slightly separated, thus originating the structure with three spots.

As far as a realistic test of multiturn extraction is concerned, it is worthwhile pointing out that the installed hardware, i.e. the two kickers used to generate the closed bump five-turn long and the other elements used to extract the beam in the CT mode, does not allow for any clean extraction of the beam trapped into the stable islands.

**MULTITURN INJECTION**

An interesting feature of the proposed multiturn extraction is that it can be easily adapted to perform multiturn injection as it was already pointed out in Ref. [7]. In fact, due to the time-reversal property of the underlying Hamiltonian dynamics, it is possible to use stable islands to inject some beam into them and to merge the beamlets by varying continuously the tune [15]. During the multiturn injection proper, it is necessary to generate a fast closed bump around the injection septum by means of fast kickers. The time duration of such a bump should equal the time duration of the number of injected turns. After that, the fast bump is collapsed to zero, and the tune variation applied. An example of multiturn injection based on this novel approach is shown in Fig. 11 for the fourth-order resonance. The evolution of the beam distribution at various stages of the injection process is shown until the final merging of the beamlets. The tune variation following the injection process is reported in Fig. 12 (left). It is worth mentioning that the numerical simulations whose results are reported here, do not take into account any interaction between the particles in the beamlets.

The final result is rather interesting, as the four beamlets are not merged into a single Gaussian beam centred in zero, but they generate a hollow beam as is shown in Fig. 12 (right). This is due to the fact that the islands’ size tends to zero in a neighbourhood of the origin. Therefore, the particles trapped inside stable islands cannot be transported towards a vanishing amplitudes. This fact is an interesting feature of the proposed method, as the hollow beam distribution naturally decreases the influence of space charge forces. Hence, this approach could be applied in all situations where space charge effects are the limiting factors at injection into a circular machine. Of course, numerical simulations will be required to confirm that the peculiar transverse phase space distribution is actually preserved under the influence of electromagnetic interaction between charged particles. In case there is no need for a hollow-beam at the end of the injection process, it is possible to change slightly the
approach. In fact, as the fourth-order resonance is stable it is possible to inject a fifth turn in the centre of phase space. The properties of the core beam can be chosen so that the final distribution, resulting from the merging of the hollow and Gaussian beams, features the required shape.

As for the multiturn extraction, the number of injected turns depends on the resonance order used for the generation of the stable islands: other schemes based on the third-order resonance have been successfully simulated [15].

CONCLUSIONS AND OUTLOOK

A novel approach to perform multiturn extraction was proposed a few years ago. Following the successful results of numerical simulations, an experimental campaign was launched in the year 2002. Presently, a single-bunch proton beam can be split in the horizontal plane with no measurable losses both at low- and high-intensity. The generated beamlets can be sufficiently separated to allow a virtually loss-free extraction. A number of issues are still under discussion before the proposed method can be considered a viable solution to the present Continuous Transfer. In particular, the fraction of trapped particles should be increased from the achieved value of 13 – 14 % to the nominal value of 20 %.

Based on the same approach, but reversing the time, a multiturn injection could be envisaged by injecting beam into the stable islands. According to the results of numerical simulations the method is working well. Indeed, it allows generating hollow beams, which might be an interesting feature whenever space charge forces are a limiting factor at injection. This point should be confirmed by numerical simulations taking into account the interaction between the particles. Furthermore, whenever is not required to have a fancy final beam distribution, the shape can be fine-tuned by injecting an additional beam at the origin of phase space. It is worthwhile stressing that the approach based on trapping particles inside stable islands of transverse phase space is, indeed, a mean of manipulating the transverse emittance in a synchrotron. Thanks to this novel technique, the spectrum of possible beam manipulations is highly increased. It is customary to impose that the transverse emittance be conserved along the injectors’ chain of a high-energy collider, with the emittance value dictated by the aperture and the required luminosity of the final machine. Indeed, by using the approach proposed here, the transverse beam emittance might be much larger than the actual required value in the low-energy part of the injectors’ chain, as the emittance could be tailored by means of multiturn extractions between the various machines. For instance, this strategy could be applied to the problem of increasing the filling efficiency of the LHC machine. In fact, the emittance generated by the PS-Booster could be much larger than the present value, the reduction being performed between PS and SPS by applying a multiturn extraction. An appealing side effect would be a much shorter filling time of both the SPS and the LHC, thus improving the availability of the CERN machines for other physics users.

ACKNOWLEDGEMENTS

We would like to thank M. E. Angoletta for support with the turn-by-turn acquisition system, M. Benedikt and M. Chanel for preparing the various beams in the PS-Booster, S. Hancock for setting up the longitudinal parameters in the PS and C. Bal, B. Dehning, J. Koopman, U. Raich for the support with the wire scanner.

REFERENCES

Abstract

This paper gives a brief history and a general description of fixed field alternating gradient (FFAG) accelerators, from the early years till the most recent proton accelerator projects.

INTRODUCTION

The concept of fixed field alternating gradient (FFAG) accelerators dates from the early 1950’s [1]. They were seen as a way to apply the principles of strong focusing and synchrotron stability and yielded high intensity machines, at a time where fixed orbit strong focusing synchrotrons eventually took over, while cyclotrons were limited to lower energies.

Only 5 FFAGs have been constructed and operated up to now:
- 3 electron machines in the 1950’s, by MURA (Midwestern Universities Research Association), which saw many crucial tasks of accelerator physics first tackled [1],
- 2 proton machines by KEK recently, in a context of determining technological progress regarding magnetism, acceleration and other beam manipulation equipments.

Nevertheless, FFAGs have regularly been proposed as an alternative solution to Linac, RCS and other cyclotron, for the production of proton beams. More recently, the neutrino factory studies triggered strong R&D activity in the field, and on the other hand the emergence of new concepts as well as modern technologies have revived the interest in the method, and pushed to re-exploring potential applications [2].

FFAGs are one of the most active fields in accelerator research today, with 9 workshops from Dec. 1999 till Oct. 2004, and again two planned for 2005.

All these aspects will be addressed in the following, briefly though. A large amount of References will however, be of some help to the interested reader for digging into the subject.

THE MURA ELECTRON FFAGS

First model, radial sector FFAG, Mark II

Work on “Mark II” began in 1955, 2 years after the invention of the concept. The machine (Fig. 1) [3] was first operated in March 1956, at the University of Michigan. The model was to be a proof of the FFAG principle, it eventually had a rich history of demonstrating experiments regarding effects of resonances, RF acceleration, beam stacking, RF KO, etc. The magnetic field is by principle fixed in time, with mid-plane form \( B(r, \theta) = B_0 \left( \frac{r}{r_0} \right)^K \mathcal{F}(\theta) \) \((K > 0\) a constant, \(r_0\) a reference radius, \(\mathcal{F}(\theta)\) an axial form factor) (see plot below), fast increasing with radius, from lower energy, larger gap, on inner orbit (where the beam is injected) to largest energy, smallest gap, on outer orbit.

The \(r^K\) shape of \(B\) is due for part to the gap size decreasing with \(r\), and for the rest to coil winding arrangement thus allowing \(K\) (and hence tunes) to be varied.

<table>
<thead>
<tr>
<th>(E_{\text{inj}} - E_{\text{max}})</th>
<th>keV</th>
<th>25 - 400</th>
</tr>
</thead>
<tbody>
<tr>
<td>orbit radius</td>
<td>m</td>
<td>0.34 - 0.50</td>
</tr>
<tr>
<td>lattice</td>
<td>(\frac{2\pi F}{f})</td>
<td></td>
</tr>
<tr>
<td>number of cells</td>
<td></td>
<td>8</td>
</tr>
<tr>
<td>fi eld index (K), tunable</td>
<td></td>
<td>(\approx 3.4)</td>
</tr>
<tr>
<td>(\nu_r / \nu_z), tunable</td>
<td></td>
<td>2.2-3 / 1-3</td>
</tr>
<tr>
<td>Magnet</td>
<td>radial sector</td>
<td></td>
</tr>
<tr>
<td>P, D sectors</td>
<td>deg</td>
<td>25.74, 10.44</td>
</tr>
<tr>
<td>gap, max.-min.</td>
<td>cm</td>
<td>6 - 4</td>
</tr>
<tr>
<td>Injection</td>
<td>continuous or pulsed</td>
<td></td>
</tr>
<tr>
<td>Acceleration</td>
<td>betatron core, at fi rst,...</td>
<td></td>
</tr>
<tr>
<td>swing</td>
<td>Gauss</td>
<td>40 - 150</td>
</tr>
<tr>
<td>rep. rate</td>
<td>(H_C)</td>
<td>a few 10's</td>
</tr>
<tr>
<td>freq. swing</td>
<td>MHz</td>
<td>10 in [35, 75]</td>
</tr>
<tr>
<td>gap voltage</td>
<td>(V)</td>
<td>50</td>
</tr>
</tbody>
</table>

The ring is built from an alternance (hence the \(\mathcal{F}(\theta)\) form factor) of positive dipoles which yield radial focusing \((\frac{d^2}{ds^2}) \frac{d\beta}{ds} > 0\) and shorter, negative dipoles which yield radial defocusing \((\frac{d^2}{ds^2}) \frac{d\beta}{ds} < 0\), thus insuring AG strong focusing. The radial dependence \(B = B_0 (r/r_0)^K\) determines the “scaling” property (also known as the “zero-chromaticity condition”): tunes are independent of the orbit (hence, of energy), closed orbits are similar wrt. ge-
metrical center (they have a scalloped shape, due to the alternating curvature). Series of basic properties ensue, like a large circumference factor $C/2\pi \rho$, momentum compaction $\alpha = 1/(1 + K)$, $\gamma r = \sqrt{1 + K}$ easily put beyond top energy, feasibility of arbitrary RF programs: no need to track $B$, and so forth. In the linear approximation the motion about a closed orbit satisfies Hill’s equations $z'' + \frac{1-n}{2} \frac{m}{r^2} z = 0$, $z'' + \frac{m}{r^2} z = 0$ with $n(s) \approx -\frac{d}{dr} \frac{dB}{dr} = -K/C$, thus amenable to regular optical treatment, working point gymnastics, defect analysis, etc. The longitudinal motion in presence of RF obeys, as in synchrotrons, $\dot{\phi} + \frac{q^2}{\cos \psi}(\sin \phi - \sin \phi_s) = 0$. The Table above gives the main parameters of Mark II.

Second model, spiral sector FFAG, Mark V

Work on “Mark V” began in 1955, a year after the invention of the concept. The machine (Fig. 2) [4] was first operated in August 1957 in the MURA Lab, Madison. Objectives were to validate the strong focusing spiral optics with its advantage of a smaller circumference, and perform beam physics, accelerator studies.

### Table

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E_{inj} - E_{max}$ (keV)</td>
<td>35 - 180</td>
</tr>
<tr>
<td>orbit radius (m)</td>
<td>0.34 - 0.52</td>
</tr>
<tr>
<td>number of sectors</td>
<td>6</td>
</tr>
<tr>
<td>fi eld index $K$, tunable</td>
<td>0.7</td>
</tr>
<tr>
<td>$\nu_1/\nu_2$, tunable</td>
<td>1.1</td>
</tr>
<tr>
<td>$\beta_1/\beta_2$ (m)</td>
<td>0.45-1.3/0.6-1.4</td>
</tr>
<tr>
<td>Magnet (spiral sector)</td>
<td>1.4/1.2</td>
</tr>
<tr>
<td>edge radius angle, $Atg(Nw)$ (deg)</td>
<td>46</td>
</tr>
<tr>
<td>$r_{min} - r_{max}$ (m)</td>
<td>0.25 - 0.61</td>
</tr>
<tr>
<td>gap, max.-min. (cm)</td>
<td>16.5 - 7</td>
</tr>
<tr>
<td>Injection</td>
<td>continuous or pulsed</td>
</tr>
<tr>
<td>Acceleration (betatron &amp; RF)</td>
<td>1000 - 0.1</td>
</tr>
<tr>
<td>RF voltage (V)</td>
<td>150</td>
</tr>
</tbody>
</table>

The idea in the spiral FFAG was to superpose a positive field on top of the alternating sign of the radial sector case, so as to always have the right curvature and hence decrease the circumference factor, which yields the “Thomas focusing” of cyclotrons. Yet by doing so the vertical focusing is weakened and needs be recovered by spiraling the poles. Appropriate field form for insuring the scaling property and constant closed orbit to spiral edge angle, is $B(r, \theta)|_{z=0} = B_0 (r/r_0)^K F \left( \frac{\ln \left( r/r_0 \right) }{w} - \theta \right)$. The axial modulation $F$ is called the “flutter”, it has the approximate form $F = 1 + f \sin(\gamma \ln \left( r/r_0 \right) /w - \theta)$. Expansion of the equations of motion around the closed orbit in the linear approximation (or as well a hard edge matrix model) yields the tunes $\nu_c \approx \sqrt{1 + K}$, $\nu_s \approx \sqrt{-K + (f/Nw)^2}/2$. The Table above gives the main parameters of Mark V.

### 50 MeV, two-way, electron FFAG

Work on the 50 MeV electron FFAG began in 1957 [5]. The machine (Fig. 3) was first operated in 1959 with two 27 MeV beams stored in opposing directions, as made possible by the radial sector optics using identical dipoles in a FODO arrangement. 51 MeV energy, one-way, was reached in 1960 after modifications in the magnets. Colliding beams, once envisaged, a hot task in the mid-50’s, need intensity, RF stacking was developed and allowed it, 10 amperes intensity was obtained that way. The Table above gives the FFAG parameters.

### Table

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E_{inj} - E_{max}$ (MeV)</td>
<td>0.1 - 50</td>
</tr>
<tr>
<td>orbit radius (m)</td>
<td>1.2 - 2.00</td>
</tr>
<tr>
<td>number of cells</td>
<td>16</td>
</tr>
<tr>
<td>$K$</td>
<td>9.25</td>
</tr>
<tr>
<td>$\nu_1/\nu_2$</td>
<td>4.42 / 2.75</td>
</tr>
<tr>
<td>Magnet (radial sector)</td>
<td></td>
</tr>
<tr>
<td>sector angle (deg)</td>
<td>6.3</td>
</tr>
<tr>
<td>peak fi eld $T$ (º)</td>
<td>0.52</td>
</tr>
<tr>
<td>gap, max.-min. (cm)</td>
<td>8.6 - 8.0</td>
</tr>
<tr>
<td>Acceleration (betatron &amp; RF)</td>
<td>20 - 23</td>
</tr>
<tr>
<td>RF swing (MHz)</td>
<td>1.3 - 3</td>
</tr>
<tr>
<td>cycle rep. rate (Hz)</td>
<td>60</td>
</tr>
</tbody>
</table>

The KEK Proton Machines

### POP

KEK POP (proof of principle) machine (Fig. 4) [6] is the first proton FFAG, first operated in 2000. Its design has strongly benefited from modern magnet computation tools and sophisticated tracking codes. The DFD lattice allows comfortable drifts, it is based on a radial sector triplet (two negative dipoles at both ends and...
Figure 4: POP FFAG, FDF dipole triplet with scaling gap, broad band accelerating cavity.

A larger, positive one in between, in a common yoke with so-called “scaling” gap shape $g_0(r/r_0)^K$ producing the radial field dependence $B_0(r/r_0)^K$. The acceleration uses high gradient, broad band “FINEMET” technologies yielding a narrow cavity (see Fig. 4) and a potential 1 kHz rep. rate [7]. Injection is on inner radius, via an electrostatic inflector, either single-turn (using a chopper in the injection line) or multi-turn (using two bump electrodes). Tunes are adjustable via the $B_F/B_D$ ratio. The Table below, col. “POP”, gives the main parameters of the POP FFAG.

### POP 150 MeV

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
<th>150 MeV Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>$E_{inj}$ - $E_{max}$ (MeV)</td>
<td>0.05 - 0.5</td>
<td>12 - 150</td>
</tr>
<tr>
<td>Orbit radius (m)</td>
<td>0.8 - 1.14</td>
<td>4.7 - 5.2</td>
</tr>
<tr>
<td>Number of cells</td>
<td>8</td>
<td>12</td>
</tr>
<tr>
<td>$K$</td>
<td>2.5</td>
<td>7.6</td>
</tr>
<tr>
<td>$\beta_r / \beta_z$ max. (m)</td>
<td>0.7 / 0.7</td>
<td>3.8 / 1.3</td>
</tr>
<tr>
<td>$\nu_r / \nu_z$</td>
<td>2.2 / 1.25</td>
<td>3.7 / 1.2</td>
</tr>
<tr>
<td>Magnet</td>
<td>DFD</td>
<td>radial sector</td>
</tr>
<tr>
<td>D, F sectors deg</td>
<td>2.8 / 14</td>
<td>3.43 / 10.24</td>
</tr>
<tr>
<td>$B_D/B_F$, max-min (T)</td>
<td>.04-.13/.14-.32</td>
<td>.3-.8 / .5-1.6</td>
</tr>
<tr>
<td>Gap, max.-min. (cm)</td>
<td>30 - 9</td>
<td>20 - 4</td>
</tr>
<tr>
<td>Swing (MHz)</td>
<td>0.6 - 1.4</td>
<td>1.5 - 4.6</td>
</tr>
<tr>
<td>Voltage p-to-p (kV)</td>
<td>1.3 - 3</td>
<td>19</td>
</tr>
<tr>
<td>Cycle time (ms)</td>
<td>1</td>
<td>4</td>
</tr>
<tr>
<td>Rep. rate (Hz)</td>
<td>$10^3$</td>
<td>250</td>
</tr>
<tr>
<td>$\dot{B}$ equivalent (T/s)</td>
<td>180</td>
<td>280</td>
</tr>
</tbody>
</table>

150 MeV proton FFAG

This second, higher energy, proton FFAG was first operated in 2003 [8]. The structure is similar to POP, it uses a 10 MeV cyclotron injector. One distinguishing feature is the return-yoke free dipole triplet (see lower-left corner in the photo, Fig. 5) which facilitates beam injection and extraction. Due to the extending fringe fields and to saturation effects, the zero-chromaticity condition is not fully satisfied, so that tunes slightly vary over the energy span. The project has various goals, as investigating applications to cancer proton therapy, accelerator driven systems, and includes R&D related to high repetition rate, fast extraction, etc. The Table above, col. “150 MeV”, gives the main parameters of the machine.

### Tracking

A remark arises from experience: the design of FFAG machines must resort to tracking, possibly using field maps [9], as early as the first order design stages, in order to access closed orbits, tunes, and other optical functions. Analytic or matrix approach can only yield approximate values of zeroth and first order parameters, only good as starting guidelines [3]. That specificity of FFAG design was already clear in the early years, where digital computation was abundantly used in field and trajectory calculations [1]. In addition, tracking is the only way one can access transverse stability limits (Fig. 6 left), amplitude or momentum detuning, 6-D acceleration (Fig. 6 right), etc.

Figure 5: KEK 150 MeV proton FFAG

Figure 6: Horizontal motion limits (left) and 12→150 MeV acceleration (right) in the KEK 150 MeV FFAG [10].
Precision 6-D tracking is of prime importance for instance when comparing muon FFAGs (see last Section) based either on “scaling” optics (strongly non-linear transverse motion) or on “non-scaling” optics (strongly non-linear longitudinal motion).

FROM 1964 TO TODAY’S R&D

After the MURA years, some activity kept going on FFAG, usually in devising alternatives to Linac or synchrotron designs in (high power) proton beam based projects, with such possible advantages as their allowing low circulating current, or lower investment cost. Let us mention for illustration, the European spallation neutron source project (ESS) [11], based on a MW range pulsed proton beam, that lead to the FFAG parameters below (col. A in the Table) [12], and the Fermilab 8 GeV proton driver project that lead to two FFAG alternative proposals, a spiral sector and a radial sector design (col. B in the Table) [13].

<table>
<thead>
<tr>
<th></th>
<th>A</th>
<th>B</th>
</tr>
</thead>
<tbody>
<tr>
<td>beam power</td>
<td>MW</td>
<td>5</td>
</tr>
<tr>
<td>top E</td>
<td>GeV</td>
<td>3</td>
</tr>
<tr>
<td>p/pulse</td>
<td>2 \times 10^{14}</td>
<td>3.6 \times 10^{12}</td>
</tr>
<tr>
<td>rep. rate</td>
<td>Hz</td>
<td>50</td>
</tr>
<tr>
<td>radius</td>
<td>m</td>
<td>140</td>
</tr>
<tr>
<td>injection E</td>
<td>MeV</td>
<td>430</td>
</tr>
<tr>
<td>DFD sectors</td>
<td>20</td>
<td></td>
</tr>
<tr>
<td>K</td>
<td>21</td>
<td></td>
</tr>
<tr>
<td>magnet width</td>
<td>m</td>
<td>2.5</td>
</tr>
<tr>
<td>RF freq./voltage kV</td>
<td>1.6-2 / 200</td>
<td>7.5 / -</td>
</tr>
</tbody>
</table>

Drawbacks in these types of tentatives were of various nature, concerning generally magnet size, insertion in an existing installation, high power beam injection in short drifts, operation costs, etc. In a general manner, large apertures that characterize scaling FFAGs entail massive magnets, radial sector optics entails large circumference.

Today’s trends

However, fixed field allows high average intensity, whereas large apertures entail large geometrical acceptance and the zero-chromaticity condition yields large momentum acceptance. As a consequence, the scaling FFAG method is still actively considered, benefiting in particular from modern technologies as stressed earlier, including as well scaling field SC magnet developments [14]. As a matter of fact, many contemporary Japan constructions were launched in this context : ADS proton driver, muon beam manipulation, protontherapy machine, high power electron beams, etc. [15], as well as the large acceptance, fast acceleration of muons in the neutrino factory [16]. As to proton beams, a recent table of comparative performance of FFAG, cyclotron and synchrotron machines can be found in Ref. [17].

On the other hand, new concepts have arisen these last years, in particular that of “non-scaling” FFAG, which contribute to their “rebirth”, this is the subject of the following.

NON-SCALING FFAGS

NuFact works have entailed a strong activity in the new field of “non-scaling” FFAGs [18, 19] presumed to bring advantages compared to classical “scaling” FFAG as involved in the Japan NuFact design, in particular in terms of lower cost in the higher energy stages of muon acceleration, and in their allowing the use of high frequency / high gradient SC RF.

“Non-scaling” optics has the large energy acceptance proper to FFAG, it was at first based on linear, combined function magnets, therefore prone to large dynamic aperture [20]. By “non-scaling” it is meant that tunes are allowed to vary in the course of acceleration (in practice, a decrease of the cell tune, due to the natural chromaticity, of about a 1/2 integer). In the muon application for instance, multi-GeV acceleration using a hundreds of cells ring means crossing “forests” of Floquet’s resonances over the few turns in the ring from injection to top energy, fast enough though, not to yield prohibitive constraints on magnet alignment and defects. Other features of “non-scaling” optics is, a better circumference factor, smaller aperture magnets compared to scaling FFAGs and in particular smaller dispersion, lower fields, the possibility of near-crest acceleration [21].

The concept has been extended to non-linear magnetic fields, with dramatic consequences on the lattice properties. Non-linear, non-scaling FFAG optics permits such design as isochronous lattice [22], allowing on-crest acceleration, as in cyclotrons. Additional sophistication in the spatial behavior of the magnetic field have also allowed designing weakly non-scaling lattices [23] in which the total tune only varies by a fraction of an integer, making this type of optics a good candidate for various applications [19], amongst which proton drivers as addressed below.

No “non-scaling” FFAG has ever been built, which motivates a recent proposal for an electron model, in the 10s MeV range [24].

These new concepts are now envisaged as an alternative to “scaling” optics in the regular fields of interest of FFAGs, as hadrontherapy, high power proton beams, etc.

An example of a proton booster application

The 1 MW upgrade of the AGS at BNL requires increasing the repetition rate to 2.5 Hz, and the number of particles to 10^{14}ppp. This imposes replacing the 1.5 GeV booster ring, the baseline scenario being based on a superconducting Linac, whereas a non-scaling FFAG appears to be a cost effective alternative.

A non-scaling optics has recently been worked out, based on an “adjusted field profile” that causes the index to be a function of the radial displacement x and of the longitudinal position in the dipoles, that is \theta = \theta(x, \theta) (Fig. 7), with the effect of cancelling the momentum dependence of the focusing strength. A consequence of this field shape is a reduced “non-scaling” : the variation of the total tune is only of the order of a fraction of an integer over

\[ n = n(x, \theta) \]
Figure 7: Field profiles vs. radial excursion at some azimuths in the F- (left graph) and D-sector (right) AFP magnets.

Figure 8: Left : total tune values during acceleration. Right : optical functions in the adjusted field profile non-scaling FDF cell.

applied with dipole triplet cells that have been shown to be advantageous, especially in the FDF configuration that yields low dispersion (Fig. 8). This type of design is believed to yield competitive technology that can allow beam performance at the level of the other accelerator architectures. A main feature is in the compactness of the magnets ensuing from the much reduced beam excursion, compared to scaling FFAG.

REFERENCES


[22] Isochronous non-scaling FFAG lattice, G. Rees, see [15]. See also, A modern answer in matter of precision tracking : stepwise ray-tracing, F. Méot, these proceedings.


OVERVIEW OF SINGLE-BEAM COHERENT INSTABILITIES IN CIRCULAR ACCELERATORS

E. Métral, CERN, Geneva, Switzerland

Abstract

Single-bunch and coupled-bunch instability mechanisms in both longitudinal and transverse planes are reviewed. Stabilization by Landau damping, linear coupling, or feedbacks are also discussed. Benchmarking with some instability codes are shown as well as several experimental results.

INTRODUCTION

Two approaches are usually used to deal with collective instabilities. One starts from the single-particle equation while the other solves the Vlasov equation, which is nothing else but an expression for the Liouville conservation of phase-space density seen by a stationary observer. In the second approach, the motion of the beam is described by a superposition of modes, rather than a collection of individual particles. The detailed methods of analysis in the two approaches are different, the particle representation is usually conveniently treated in the time domain, while in the mode representation the frequency domain is more convenient, but in principle they necessarily give the same final results. The advantage of the mode representation is that it offers a formalism that can be used systematically to treat the instability problem.

The first formalism was used by Courant and Sessler to describe the transverse coupled-bunch instabilities [1]. In most accelerators, the RF acceleration mechanism generates azimuthal non-uniformity of particle density and consequently the work of Laslett, Neil and Sessler for continuous beams [2] is not applicable in the case of bunched beams. Courant and Sessler studied the case of rigid (point-like) bunches, i.e. bunches oscillating as rigid units, and they showed that the transverse electromagnetic coupling of bunches of particles with each other can lead (due to the effect of imperfectly conducting vacuum chamber walls) to a coherent instability. The physical basis of the instability is that in a resistive vacuum tank, fields due to a particle decay only very slowly in time after the particle has left (long-range interaction). The decay can be so slow that when a bunch returns after one (or more) revolutions it is subject to its own residual wake field which, depending upon its phase relative to the wake field, can lead to damped or anti-damped transverse motion. For $M$ equi-populated equi-spaced bunches, $M$ coupled-bunch mode numbers exist ($n = 0, 1, \ldots, M - 1$), characterized by the integer number of waves of the coherent motion around the ring. Therefore the coupled-bunch mode number resembles the azimuthal mode number for coasting beams, except that for coasting beams there is an infinite number of modes. The bunch-to-bunch phase shift $\Delta \phi$ is related to the coupled-bunch mode number $n$ by $\Delta \phi = 2\pi n / M$. Pellegrini [3] and, independently, Sands [4,5] then showed that short-range wake fields (i.e. fields that provide an interaction between the particles of a bunch but have a negligible effect on subsequent passages of the bunch or of other bunches in the beam) together with the internal circulation of the particles in a bunch can cause internal coherent modes within the bunch to become unstable. The important point here is that the betatron phase advance per unit of time (or betatron frequency) of a particle depends on its instantaneous momentum deviation (from the ideal momentum) in first order through the chromaticity and the slippage factor. Considering a non-zero chromaticity couples the betatron and synchrotron motions, since the betatron frequency varies around a synchrotron orbit. The betatron phase varies linearly along the bunch (from the head) and attains its maximum value at the tail. The total betatron phase shift between head and tail is the physical origin of the head tail instability. The head and the tail of the bunch oscillate therefore with a phase difference, which reduces to rigid-bunch oscillations only in the limit of zero chromaticity. A new (within-bunch) mode number $m = \ldots, -1, 0, 1, \ldots$ also called head-tail mode number, was introduced. This mode describes the number of betatron wavelengths (with sign) per synchrotron period. It can be obtained by superimposing several traces of the directly observable average displacement along the bunch at a particular pick-up. The number of nodes is the mode number $|m|$. The work of Courant and Sessler, or Pellegrini and Sands, was done for particular impedances and oscillation modes. Using the Vlasov formalism, Sacherer unified the two previous approaches, introducing a third mode number $q = \ldots, -1, 0, 1, \ldots$, called radial mode number, which comes from the distribution of synchrotron oscillation amplitudes [6,7]. The advantage of this formalism is that it is valid for generic impedances and any high order head-tail modes. This approach starts from a distribution of particles (split into two different parts, a stationary distribution and a perturbation), on which Liouville theorem is applied. After linearization of the Vlasov equation, one ends up with Sacherer’s integral equation or Laclare’s eigenvalue problem to be solved [7]. Because there are two degrees of freedom (phase and amplitude), the general solution is a twofold infinity of coherent modes of oscillation $(m, q = \ldots, -1, 0, 1, \ldots)$. At sufficiently low intensity, only...
the most coherent mode \( q = m \) (largest value for the coherent tune shift) is generally considered, leading to the classical Sacherer’s formulae in both transverse and longitudinal planes. For protons a parabolic density distribution is generally assumed, which is a reasonable approximation at relatively low energy, and the corresponding oscillation modes are sinusoidal. For electrons, the distribution is usually Gaussian, and the oscillation modes are described in this case by Hermite polynomials. In reality, the oscillation modes depend both on the distribution function and the impedance, and can only be found numerically by solving the (infinite) eigenvalue problem. However, the mode frequencies are not very sensitive to the accuracy of the eigenfunctions. Similar results are obtained for the longitudinal plane.

**TRANSVERSE**

**Low Intensity**

At low intensity (i.e. below the intensity threshold given in the next section), the standing-wave patterns (head-tail modes) are treated independently. This leads to instabilities where the head and the tail of the bunch exchange their roles (due to synchrotron oscillation) several times during the rise-time of the instability. The exchange their roles (due to synchrotron oscillation) instabilities where the head and the tail of the bunch (head-tail modes) are treated independently. This leads to given in the next section), the standing-wave patterns (sinusoidal modes for parabolic bunches). For zero chromaticity, the power spectrum of mode \( n = m \) is peaked near \( \omega = (|q|+1)\tau_\beta \) and extends \( \pm 2\pi/\tau_\beta \) (rad/s). The theoretical average displacement along the bunch at a particular pick-up is shown in Fig. 2 in an example case.

Here, \( j = \sqrt{-1} \) is the imaginary unit, \( e \) is the elementary charge, \( \beta \) and \( \gamma \) are the relativistic velocity and mass factors, \( I_p = N_p e \Omega_0 / (2\pi) \) is the current in one bunch, \( m_0 \) is the proton rest mass, \( \Omega_0, \Omega_0, \Omega_0 \) are the unperturbed betatron tunes, \( L = \beta c \tau_\beta \) is the total (4\( \pi \)) bunch length (in metres), \( Z_{x,y} \) are the transverse coupling impedances, \( \omega_{k}^{x,y} = \left( k + \Omega_{x,y,0} \right) / \Omega_0 + m_0 \) with \( -\infty < k \leq +\infty \) for a single-bunch beam, and \( k = n_{x,y} + k' M \) with \( -\infty < k' \leq +\infty \) for a multi-bunch beam (\( n_{x,y} \) are the coupled-bunch mode numbers), \( \omega_{k}^{x,y} = \left( \xi_{x,y} / \eta \right) \Omega_{x,y,0} \Omega_0 \) are the transverse chromatic frequencies, with \( \xi_{x,y} = \left( dQ_{x,y} / dp \right) \left( \eta_0 / \Omega_{x,y,0} \right) \) the chromaticities, and \( \eta = \gamma / \beta - \gamma^2 \) is the slippage factor. The bunch spectrum \( \delta_h \), represented in Fig. 1 for the first three diagonal modes, describes the cross-power densities of the \( m \)th and \( q \)th line-density modes (sinusoidal modes for parabolic bunches). For zero chromaticity, the power spectrum of mode \( |m| \) is peaked near \( \omega = (|q|+1)\tau_\beta \) and extends \( \pm 2\pi/\tau_\beta \) (rad/s). The theoretical average displacement along the bunch at a particular pick-up is shown in Fig. 2 in an example case.

\[
\Delta \nu_{m,q} = \left( |m|+1 \right)^{-1} \frac{j e \beta I_b}{2 m_0 \gamma \Omega_{x,y,0} \Omega_0 L} \left( Z_{x,y}^{m,q} \right), \tag{1}
\]

with
\[
\left( Z_{x,y}^{m,q} \right) = \sum_{k=-\infty}^{k=+\infty} \left( Z_{x,y}^{m,q} \right) h_{m,q} \left( \omega_{k}^{x,y} - \omega_{y,z} \right), \tag{2}
\]

\[
h_{m,q}(\omega) = \frac{\varepsilon^2}{\pi} \left( |m|\times |q|+1 \right) \times \sum_{k=-\infty}^{k=+\infty} \left( \omega_{k}^{x,y} - \omega_{y,z} \right) \times \left( (\alpha \tau_b / \pi)^2 - (|m|+1)^2 \right), \tag{3}
\]

\[
P_{m}^{even} = (-1)^{\frac{1}{2} \times \left( |m|+1 \right)} \times \sin \left( \omega \tau_b / 2 \right), \tag{4}
\]

\[
P_{m}^{odd} = \frac{(-1)^{|m|+1}}{2} \times \sin \left( \omega \tau_b / 2 \right), \tag{5}
\]

\[
P_{m}^{even} = \frac{(-1)^{|m|+1}}{2} \times \sin \left( \omega \tau_b / 2 \right), \tag{6}
\]

\[
P_{m}^{odd} = (-1)^{\frac{1}{2} \times \left( |m|+1 \right)} \times \sin \left( \omega \tau_b / 2 \right). \tag{7}
\]

![Figure 1: Power spectrum for the first three diagonal modes, \( m = 0 \) (purple), \( |m| = 0 \) (orange), and \( |m| = 2 \) (green).](image1.png)

![Figure 2: Five superimposed AR-signals at a Beam Position Monitor (BPM) from theory, for \( m = 0 \) (upper), and \( |m| = 1 \) (lower), in an example case where \( \xi \neq 0 \). The number of nodes is equal to the modulus of the head-tail mode number \( |m| \). A particular turn is shown in red.](image2.png)
Two experiments performed at the CERN Proton Synchrotron (PS) are reviewed in the following and compared to the theoretical expectations. In the first, a multi-bunch proton beam was observed to be unstable in 1997, suffering from a coupled-bunch instability [8]. From Sacherer’s theory, the predicted most critical head-tail mode was $|m| = 1$ (see Fig. 3), which was confirmed by measurements (see Fig. 4). In the second experiment, a single-bunch instability was observed with an LHC-type beam in the PS in 1999 [9]. Here also the predictions were in good agreement with the observations (see Figs. 5 and 6). Figure 7 exhibits different unstable modes ($|m| = 4, 5, 7, 8, 10$) in the horizontal plane, in good agreement with Sacherer’s theory, which have been obtained by tuning the chromaticity. No stabilizing values of chromaticity were found.

Figure 3: (Upper) classical resistive-wall impedance with CERN PS parameters (“thick-wall” case), and (lower) power spectrum for the first five head-tail modes.

Figure 4: $\Delta R$-signal from a radial BPM during 20 consecutive turns. As predicted from Sacherer’s theory, the head-tail mode $|m| = 1$ is observed.

Figure 5: Growth rate vs. the head-tail mode number. The head-tail mode number $m = 6$ is predicted to be the most unstable.

Figure 6: $\Delta R$-signal from a radial BPM during 20 consecutive turns. As predicted from Sacherer’s theory, the head-tail mode $m = 6$ is observed.

Figure 7: $\Delta R$-signals from a radial BPM during 20 consecutive turns, obtained by tuning the chromaticity. Time scale: 20 ns/div.
High Intensity

As the bunch intensity increases, the different head-tail modes can no longer be treated separately. In this regime, the wake fields couple the modes together and a wave pattern travelling along the bunch is created: this is the Transverse Mode Coupling Instability (TMCI). The TMCI for circular accelerators has been first described by Kohaupt [10] in terms of coupling of Sacherer’s head-tail modes. This extended to the transverse motion, the theory proposed by Sacherer [11] to explain the longitudinal microwave instability through coupling of the longitudinal coherent bunch modes. In linear accelerators, the Beam Break-Up (BBU) theory has been developed to explain the observed beam emittance growths and the transverse instabilities [12,13]. It has been known for some time, using a two-particle model, that the TMCI is the manifestation in synchrotrons of the BBU mechanism observed in linacs [14,15]. The only difference comes from the synchrotron oscillation, which stabilizes the beam in synchrotrons below a threshold intensity by swapping the head and the tail continuously. This effect disappears close to transition energy, or more generally when the instability rise-time is much faster than the synchrotron period. In this case, it is usually said that the concept of head-tail modes loses its meaning and that it is appropriate to use the BBU formalism to describe the interaction between the beam and its surroundings [14].

Other formalisms have also been developed to describe the instability when the rise-time is faster than the synchrotron period [16,17]. Therefore, several analytical formalisms exist for fast (compared to the synchrotron period) instabilities, but the same formula is in fact obtained (within a factor smaller than 2) from the five, seemingly diverse, formalisms in the case of a Broad-Band (BB) resonator impedance \( Z^BB_s \) [18]: (i) Coasting-beam approach with peak values [19], (ii) Fast blow-up [16], (iii) Beam break-up (for 0 chromaticity) [15], (iv) Post head-tail [17], and (v) TMCI with 2 modes in the “long-bunch” regime (for 0 chromaticity) [20]. Two regimes are indeed possible for the TMCI according to whether the total (4\( \pi \)) bunch size is inversely proportional to the value of the (vertical) BB resonator impedance \( Z^BB_y \) or vs. the resonance frequency \( f_y \) [19].

The quasi coasting-beam approach using the peak values of bunch current and momentum spread as input for the coasting-beam formula yields the following threshold for the number of protons per bunch [19]

\[
N^b_{th} = 32 \frac{\sqrt{2}}{3} \frac{Q_{v0}|\eta|\epsilon_j}{e \beta^2 c} \times \frac{f_r}{f_y} \times \left[ 1 + \frac{f_z}{f_r} \right],
\]

Equation (8) has been compared to two kinds of codes [21]. The first is MOSES [22], which is a program computing the coherent bunched-beam modes. The second is HEADTAIL [23], which is a code simulating single-bunch phenomena. These two codes have been compared to the coasting-beam formula with peak values on the following example, using the CERN Super Proton Synchrotron (SPS) parameters. A BB impedance is assumed, with \( Q_v = 1, R_s = 20 \text{ M\text{\Omega}/m} \), and a variable resonance frequency \( f_z \). The numerical values used for the CERN SPS machine are the vertical tune \( Q^y_{v0} = 26.13 \), the beam momentum \( p = 26 \text{ GeV}/c \), the slippage factor \( \eta = 6.1797 \times 10^{-4} \), the revolution frequency \( f_o = 43347.3 \text{ Hz} \), and the transverse rms beam sizes \( \sigma_y = \sigma_x = 1.2 \text{ mm} \). Three scans have been made and the results are depicted in Figs. 9 to 11. It is already known from Eq. (8) and the two codes that the intensity threshold is inversely proportional to the value of the impedance. The first scan is vs. the resonance frequency of the BB impedance, the second is vs. the bunch

![Figure 8: (Upper) power spectra for a short \((\tau_b = 0.5/f_r)\) bunch, and real and imaginary parts of the driving (vertical) BB impedance, and (lower) TMCI intensity threshold near \(2f_r, \tau_b = 1\).](image)

\[
\]
longitudinal emittance, and the third one is vs. the chromaticity.

It is seen from Eq. (8) and Figs. 9 to 11, that concerning machine parameters, the intensity threshold is increased by increasing the modulus of the slippage factor, and/or the ratio between the resonance frequency and the peak value of the resonator impedance. Concerning beam parameters, the intensity threshold is increased by increasing the longitudinal emittance, and/or the chromaticity. The first method is used in the CERN PS to avoid the fast instability at transition with high-intensity bunches (see Fig. 12) [24]. The second has been used at ESRF [17], and also in the CERN SPS with a high-intensity single-bunch beam of low longitudinal emittance (see Fig. 13) [21]. Note that as it is the longitudinal emittance which matters in Eq. (8), the Potential-Well Distortion (PWD) should have no effect on the threshold intensity for protons, as the longitudinal emittance is supposed to be conserved in this mechanism.
Figure 13: Fast instability observed in 2003 in the CERN SPS at injection ($p = 26$ GeV/c) with a high-intensity single-bunch beam ($N_b \approx 1.2 \times 10^{11}$ p/b) of low longitudinal emittance ($\epsilon_l \approx 0.2$ eVms). The synchrotron period is $T_s \approx 7.1$ ms. The bunch, which is unstable for a vertical chromaticity close to zero (upper), is stabilized by increasing the chromaticity (lower).

\begin{align*}
\xi_{y} = 0.8
\end{align*}

**LONGITUDINAL**

**Low Intensity**

The same formalism as in the transverse plane can be used. An additional complication comes here from the PWD induced by the imaginary part of the longitudinal coupling impedance, which has to be taken into account and which makes the synchrotron frequency, the bunch length and the momentum spread depend on the bunch intensity. For any intensity (below the intensity threshold discussed in the next section) the bunch length is deduced from emittance (momentum spread) conservation for protons (leptons). In addition, there is also a synchronous phase shift, which is usually a small effect, due to the real part of the longitudinal coupling impedance. These two effects apply to the stationary distribution. Taking into account these effects, a new stationary distribution is defined. Around the new fixed point, the same method as in the transverse plane can be used. A perturbation is applied, and the longitudinal bunched–beam coherent modes are deduced from the linearized Vlasov equation. The complex longitudinal coherent synchrotron frequency shift of bunched-beam modes $\Delta\omega_{m,q} = \omega_{mq} - m\omega_s$ is given by Sacherer’s formula \[11\]

\begin{align*}
\Delta\omega_{m,q}^j = \frac{m}{m+1} \times \sum_{p=-\infty}^{p=\infty} \frac{Z_i(p)}{\xi_{m,q}} h_{m,q}^\prime \left( \omega_p^l \right) \frac{p}{\xi_{m,q}^\prime}.
\end{align*}

with

\begin{align*}
\frac{Z_i(p)}{\xi_{m,q}} = \sum_{p=-\infty}^{p=\infty} \frac{Z_i(p)}{\xi_{m,q}} h_{m,q}^\prime \left( \omega_p^l \right) \frac{p}{\xi_{m,q}^\prime}.
\end{align*}

Here, $m = ..., -1, 0, 1, ...$ is the longitudinal bunched-beam mode number, $\omega_s = \omega_s(0) + \Delta\omega_s^j$ is the synchrotron angular frequency taking into account the PWD (the unperturbed synchrotron angular frequency is $\omega_s(0) = 2\pi f_s$), $B = \tau_0 f_0$ is the bunching factor taking into account the PWD (the unperturbed total bunch length is $\tau_0$). $\hat{V}_T$ is the total (effective) peak voltage taking into account the PWD (the peak RF voltage is $V_{RF}$), $h$ is the harmonic number, $\phi_s$ is the RF phase of the synchronous particle ($\cos\phi_s > 0$ below transition and $\cos\phi_s < 0$ above) taking into account the PWD (the unperturbed synchronous phase is $\phi_s(0)$), $Z_i$ is the longitudinal coupling impedance, $\omega_p^l = p\Omega_0 + m\omega_s$ with $-\infty \leq p \leq +\infty$ for a single-bunch beam, and $p = n_l + p' M$ with $-\infty \leq p' \leq +\infty$ for a multi-bunch beam ($n_l$ is the longitudinal coupled-bunch mode number, i.e. the number of waves of coherent motion per revolution). The intra-bunch mode number $m$ is the number of periods of phase space density modulation per synchrotron period in the longitudinal plane. The line density is the projection of the phase space distribution on the time axis. The observed pattern for a given bunch oscillates with $m$ times the synchrotron frequency, and it has $|m|$ nodes along the bunch.

In the absence of synchrotron frequency spread and for sufficiently low intensity, the following result is obtained for the dipole mode in the presence of an impedance with a constant imaginary part (BB + space charge)

\begin{align*}
\text{Re}(\omega_{j+1}) - \omega_s = \text{Re}(\omega_{j+1}) - \omega_s + \Delta\omega_s^j = \text{Re}(\Delta\omega_s^j) + \Delta\omega_s^j = 0.
\end{align*}

This is why it is often said that the coherent synchrotron frequency of the dipole mode does not move, and that $\text{Re}(\omega_{j+1}) \approx \omega_s$ \[7\]. However, as already known, this result is valid only in the absence of synchrotron frequency spread (see Fig. 29) and for sufficiently low intensity. Similarly, the following result is obtained for the quadrupole mode,

\begin{align*}
\text{Re}(\omega_{j+2}) - 2\omega_s = \text{Re}(\omega_{j+2}) - 2\omega_s + \Delta\omega_s^j = \text{Re}(\Delta\omega_s^j) + 2\Delta\omega_s^j \approx 2\Delta\omega_s^j / 2.
\end{align*}

The coherent synchrotron frequency of the quadrupole mode is given by $\text{Re}(\omega_{j+2}) \approx 2\omega_s + \Delta\omega_s^j / 2$ \[25\]. In conclusion, as far as frequencies are concerned, coherent and incoherent effects subtract.
**High Intensity**

In the longitudinal plane, the microwave instability for coasting beams is well understood [19,26,27,28]. It leads to a stability diagram, which is a graphical representation of the solution of the dispersion relation depicting curves of constant growth rates, and especially a threshold contour in the complex plane of the driving impedance. When the real part of the driving impedance is much greater than the modulus of the imaginary part, a simple approximation, known as the Keil-Schnell (or circle) stability criterion, may be used to estimate the threshold curve [27]. For bunched beams, it has been proposed by Boussard [29] to use the coasting-beam formalism with local values of bunch current and momentum spread. This approximation was expected to be valid in the case of instability rise-times shorter than the synchrotron period, and wavelengths of the driving wake field much shorter than the bunch length. This empirical rule is widely used for estimations of the tolerable impedances in the design of new accelerators. A first approach to explain this instability, without coasting-beam approximations, has been suggested by Sacherer through Longitudinal Mode-Coupling (LMC) [11]. The equivalence between LMC and microwave instabilities has been pointed out by Sacherer [11] and Laclare [7] in the case of BB driving resonator impedances, neglecting the PWD. The complete theory describing the microwave instability for bunched beams is still under development [28,30]. Using the mode-coupling formalism for the case of a bunch interacting with a BB resonator impedance, and whose frequency, i.e. the ratio between the momentum spread and the bunch length, a new formula has been derived taking into account the PWD due to both space-charge and BB resonator impedances, and is given by [18]

\[
\frac{Z_{BB}^l}{\Delta l} \times \left[ 1 - \text{Sgn} (\eta) \times \frac{3}{4} \left( \frac{Z_{SC}^l}{\Delta l} - 1 \right) \right]^{1/4} \leq \frac{(E/e) \beta^2 |\eta| \left( \frac{\Delta p}{p_0} \right)^2}{I_{p0}}. 
\]

(13)

Here, \( Z_{BB}^l / \Delta l \) and \( Z_{SC}^l / \Delta l \) are the peak values of the BB and space-charge longitudinal coupling impedances \( Z_{BB}^l(p)/p = R_l(\omega_l/\omega) / [j1 - jQ_l(\omega_l/\omega - \omega/\omega)] \) and \( Z_{SC}^l(p)/p = -jZ_{BB}^l / [1 + 2\ln(b/a)/(2\beta^2)] \), with \( p = \omega/\omega_l \), \( R_l \) the shunt impedance (in \( \Omega \)), \( Z_0 = 377\Omega \) the free space impedance, \( a \) and \( b \) the average beam and pipe effective radii. \( \text{Sgn} (\eta) \) denotes the sign of \( \eta \) (it is - below transition and + above), \( I_{p0} = 3 e N_b / (2 \gamma_{\text{FWHH}}) \) is the bunch peak current (without PWD) considering a parabolic line density, and \( (\Delta p/p_0)_{\text{FWHH}} \) is the full width at half height of the relative momentum spread (without PWD).

It is seen from Eq. (13) that concerning machine parameters, the threshold is increased by increasing the modulus of the slippage factor. Concerning beam parameters, the threshold is increased by increasing the energy and/or the bunch length and/or the momentum spread. Here, as opposed to the transverse case, the momentum spread is more efficient than the bunch length: for a given longitudinal emittance, short bunches are more stable than longer ones.

The stability diagram derived from Eq. (13) is shown in Fig. 14. Below transition, the space-charge impedance has a destabilizing effect but, even if the space-charge impedance is much bigger than the BB one, the effect on the threshold is rather small due to the exponent \( 1/4 \) in Eq. (13). Above transition, the space-charge impedance has a stabilizing effect, as it increases the synchrotron frequency, i.e. the ratio between the momentum spread and the bunch length.

![Stability diagram](image)

**Figure 14:** Stability diagram for the LMC instability below and above transition respectively for a proton bunch. The Keil-Schnell circle is represented by the dashed curve.

For a lepton bunch, the stability criterion writes [31]

\[
I_{p0} \leq F \times I_{p0}^{\text{KSB}}, 
\]

(14)

with

\[
I_{p0}^{\text{KSB}} = \frac{1}{\ln 2} \times \frac{(E/e) \alpha p}{(\Delta p/p_0)_{FWHH}^2} \times \left( \frac{\Delta p}{p_0} \right)^2, 
\]

(15)

where \( I_{p0}^{\text{KSB}} \) is the peak intensity threshold (without PWD) from the Keil-Schnell-Boussard criterion for Gaussian bunches [32], \( \alpha_p = \gamma_p^2 \) is the momentum...
compaction factor, and $F$ is plotted in Fig. 15. For a sufficiently long bunch, the intensity threshold is found to be ~2 times larger than from the Keil-Schnell-Boussard criterion. When the bunch length gets smaller, the threshold intensity increases. It is ~5 times larger than from the Keil-Schnell-Boussard criterion when $10^{\tau_{brf}} \approx 1$, which seems in good agreement with the (non) observations made in Ref. [33]. This may explain why the classical instability threshold has been exceeded in some lepton machines.

Experimentally, the most evident signature of the LMC instability is the intensity-dependent longitudinal beam emittance blow-up to remain just below threshold. A typical picture, obtained with a CERN PS beam during debunching at 25 GeV total energy, is shown in Fig. 16 [9].

**Figure 15:** Plot of the factor $F$ (for the LMC instability of a lepton bunch) vs. $f_s \tau_{brf}$.

**Figure 16:** Longitudinal Schottky scan spectrogram during debunching. Time goes from top to bottom. Total time window is ~200 ms. In the first 100 ms the beam is still bunched by the RF voltage, which is adiabatically decreased and then switched OFF. During the debunching there is a momentum blow-up. The last “transient” is produced by the fast extraction process.

### STABILIZATION METHODS FOR THE LOW-INTENSITY CASES

#### Transverse Landau Damping

From octupoles

Considering the case of a beam having the same normalized rms beam size $\sigma = \sqrt{\varepsilon}$ in both transverse planes, the Landau damping mechanism from octupoles of coherent instabilities, e.g. in the horizontal plane, is discussed from the following dispersion relation [34,35]

$$1 = -\Delta Q_{coh} \int_0^\infty \int_0^\infty J_x \frac{df(J_x,J_y)}{dJ_x} \frac{J_y}{Q_x - Q_y(J_x,J_y) - m Q_x}$$

with $Q_x(J_x,J_y) = Q_0 + a_0 J_x + b_0 J_y$. (17)

Here, $Q_x$ is the coherent betatron tune to be determined, $J_{x,y}$ are the action variables in the horizontal and vertical plane respectively, with $f(J_x,J_y)$ the distribution function, $\Delta Q_{coh}$ is the horizontal coherent tune shift, $Q_y(J_x,J_y)$ is the horizontal tune in the presence of octupoles, $m$ is the head-tail mode number, and $Q_x$ is the small-amplitude synchrotron tune (the longitudinal spread is neglected).

The $n$th order distribution function is assumed to be

$$f(J_x,J_y) = a \left(1 - \frac{J_x + J_y}{b}\right)^n,$$

where $a$ and $b$ are constants to be determined by normalization. The normalization of the distribution function to unity gives

$$\int_{J_x=0}^b \int_{J_y=0}^{b-J_x} dJ_x \int_{J_x=0}^b dJ_y \ f(J_x,J_y) = \frac{a b^2}{(n+1)(n+2)} = 1.$$ (19)

The average of the action variable is equal to the emittance ($\langle J_x \rangle = \varepsilon$), which gives

$$\int_{J_x=0}^b \int_{J_y=0}^{b-J_x} dJ_x \int_{J_y=0}^b dJ_y \ f(J_x,J_y) = \frac{a b^3}{(n+1)(n+2)(n+3)} = \varepsilon.$$ (20)

It can be deduced from Eqs. (19) and (20) that

$$b = (n+3) \varepsilon, \quad a = \frac{(n+1)(n+2)}{b^2}.$$ (21)
The transverse beam profile, e.g. in the horizontal plane, is given by
\[
g(x) = \frac{ab}{2\pi(n+1)} \int \frac{x^2 + p^2}{b^2 - x^2} \left(1 - \frac{x^2 + p^2}{2b}\right)^{n+1} dp_x
\]
\[
= \frac{2(n+2)[n+1]}{\pi^{2b}2(2n+3)!} \left(1 - \frac{x^2}{2b}\right)^{n+3/2}.
\] (22)

As can be seen from Eq. (22), the profile extends up to \(\sqrt{2b}\). In the case of a beam profile extending up to 6\(\sigma\) (as it should be the case in the LHC due to the collimators setting) the condition \(\sqrt{2b} = 6\sigma\) has to be satisfied, i.e. using Eq. (21), the distribution of order \(n = 15\) has to be considered.

The Gaussian distribution function is given by [35]
\[
f(J_x, J_y) = \frac{1}{\sqrt{\pi^2} \sigma} e^{-\frac{(J_x + J_y)^2}{\sigma^2}}.
\] (23)

The corresponding transverse beam profile is given by
\[
g(x) = \frac{1}{2\sqrt{\pi \sigma}} e^{-\frac{x^2}{2\sigma^2}}.
\] (24)

The transverse beam profiles for \(n = 2\), \(n = 15\), and for the Gaussian distribution are plotted in Fig. 17. As can be seen, the 15th order distribution function is very close to the Gaussian, which is not surprising as the closed expression for the nth order distribution is, from Eqs. (18) and (21),
\[
f(J_x, J_y) = \left(\frac{n+1}{n+3}\right)^2 e^{-\frac{(n+1)}{(n+3)}(J_x + J_y)^2}.
\] (25)

This expression tends to Eq. (23), i.e. the Gaussian distribution function, when \(n\) tends to infinity. This can be easily found by taking the logarithm of Eq. (25) and expanding it. A zoom of the tails of the transverse beam profiles is shown in Fig. 18. It shows that the tails of both 15th order and Gaussian distributions extend further than the quasi-parabolic \((n = 2)\) distribution by more than half a \(\sigma\), while the tails of the 15th order distribution remain below that of the Gaussian distribution.

For the nth order distribution function, the dispersion relation of Eq. (16) can be re-written as
\[
\Delta Q_{coh}^n = -a_0 e K^{-1}(c, p),
\] (30)

with
\[
K(c, p) = \int_{J_{x,0}}^{\infty} dJ_x \int_{J_{y,0}}^{\infty} dJ_y \frac{J_x e^{-(J_x + J_y)}}{p + J_x + cJ_y},
\] (31)

\[
p = \frac{Q_c - Q_0 - mQ_x}{-a_0 e}.
\] (32)

Equation (31) can be solved analytically and is given by [35]
\[
K(c, p) = \frac{1 - e^{-c(p + c - p)} e^{cE_z(p + cE_z)}}{1 - cE_zE_z},
\] (33)

where
\[
E_z(t) = \int_{t-\tau}^{t-\tau} e^{-t} dt
\] (34)

is the exponential integral function.

The l.h.s of Eq. (26) or (30) contains information about the beam intensity and the impedance. The r.h.s contains information about the beam frequency spectrum. Calculation of the l.h.s is straightforward. For a given impedance, one only needs to calculate the complex mode frequency shift, in the absence of Landau damping. Without frequency spread, the condition for the beam to be stable is thus simply \(\text{Im}(\Delta Q_{coh}^n) \geq 0\) (oscillations of
the form $e^{i\omega t}$ are considered. Once its l.h.s is obtained, Eq. (26) or (30) can be used to determine the coherent betatron tune $Q_c$ in the presence of Landau damping when the beam is at the edge of instability (i.e. $Q_c$ real). However, the exact value of $Q_c$ is not a very useful piece of information. The more useful question to ask is under what conditions the beam becomes unstable regardless of the exact value of $Q_c$ under these conditions, and Eq. (26) or (30) can be used in a reversed manner to address this question. To do so, one considers the real parameter $Q_c - Q_0 - mQ_\perp$ (stability limit) and observes the locus traced out in the complex plane by the r.h.s of Eq. (26) or (30), as $Q_c - Q_0 - mQ_\perp$ is scanned from $-\infty$ to $+\infty$. This locus defines a “stability boundary diagram”. The l.h.s of Eq. (26) or (30), a complex quantity, is then plotted in this plane as a single point. If this point lies on the locus, it means the solution of $Q_c$ for Eq. (26) or (30) is real, and this $Q_c - Q_0 - mQ_\perp$ is such that the beam is just at the edge of instability. If it lies on the inside of the locus (the side which contains the origin), the beam is stable. If it lies on the outside of the locus, the beam is unstable. The stability diagrams for the 2\textsuperscript{nd} order, 15\textsuperscript{th} order and Gaussian distribution functions are plotted in Fig. 19 for the case of the LHC at top energy (7 TeV) with maximum available octupole strength ($\varepsilon = 0.5$ nm, $|\eta_0| = 270440$ and $c = -0.65$).

![Stability diagrams](image)

**Figure 19:** Stability diagrams (positive and negative detunings $a_0$) for the LHC at top energy (7 TeV) with maximum available octupole strength, for the 2\textsuperscript{nd} order (dashed curves), the 15\textsuperscript{th} order (full curves), and the Gaussian (dotted curves) distribution.

The case of a distribution extending up to $6\sigma$ (as the 15\textsuperscript{th} order distribution) but with more populated tails than the Gaussian distribution has also been considered and revealed a significant enhancement of the stable region compared to the Gaussian case [34]. This may be the case in reality in proton machines due to diffusive mechanisms. However, as already mentioned in Ref. [37], the presence or not of the high-amplitude tails in the distribution can substantially affect the amount of Landau damping. These stability diagrams should therefore be used with great care for beam stability analyses or predictions in real machines.

Furthermore, it is important to remind that Landau damping of coherent instabilities and maximization of the dynamic aperture are partly conflicting requirements. On the one hand, a spread of the betatron frequencies is needed for the stability of the beam coherent motion, which requires nonlinearities to be effective at small amplitude. On the other hand, the nonlinearities of the lattice must be minimized at large amplitude to guarantee the stability of the single particle motion. A trade-off between Landau damping and dynamic aperture is therefore necessary.

**From both octupoles and space-charge nonlinearities**

The influence of space-charge nonlinearities on the Landau damping mechanism of transverse coherent instabilities has first been studied by Möhl and Schönauer for coasting and rigid bunched beams [38]. Later Möhl extended these results to head-tail modes in bunched beams.
beams [39]. The basic results of these studies are that in the absence of external (octupolar) nonlinearities, the space-charge nonlinearities have no effect on beam stability, as the incoherent space-charge tune spread moves with the beam. When octupoles are added, the incoherent space-charge tune spread is “mixed-in”, and in this case the octupole strength required for stabilization can depend strongly on the sign of the excitation current of the lenses.

Considering the case of a quasi-parabolic distribution function having the same normalized rms beam size $\sigma = \sqrt{\varepsilon}$ in both transverse planes and taking into account both octupoles and nonlinear space-charge forces, the Landau damping mechanism of coherent instabilities, e.g. in the horizontal plane, is discussed from the following dispersion relation [40]

$$1 = -\int \frac{dJ_x}{J_{x0}} \int \frac{dJ_y}{J_{y0}} J_x J_y \left( \frac{\partial Q_0}{\partial J_x} \frac{\partial Q_0}{\partial J_y} \right) \frac{\Delta Q_{coh} - \Delta Q_{noncoh} (J_x, J_y)}{Q_0 - Q (J_x, J_y) - m Q_0} ,$$

with

$$f(J_x, J_y) = \frac{12}{25 \varepsilon^2} \left( 1 - \frac{J_x + J_y}{5 \varepsilon} \right)^2 ,$$

$$Q_0 (J_x, J_y) = Q_0 (J_x, J_y) + \Delta Q_{noncoh} (J_x, J_y) .$$

Here, $Q_{coh} (J_x, J_y)$ is the horizontal tune in the presence of octupoles but in the absence of space-charge, given by (see Eq. (17))

$$Q_{coh} (J_x, J_y) = Q_{coh0} + a J_x + b J_y ,$$

with [41]

$$a = \frac{3}{8 \pi} \int \beta_x^2 \frac{Q_x}{B \rho} ds \quad \text{and} \quad b = -\frac{3}{8 \pi} \int 2 \beta_x \beta_y \frac{Q_y}{B \rho} ds .$$

Here, $Q_i$ is the octupolar strength, $B \rho$ the beam rigidity, and $\beta_x, \beta_y$ the horizontal and vertical betatron functions. Note that, as expected, in the absence of nonlinear space-charge forces, the dispersion relation of Eq. (16) is recovered. Furthermore, it should be noticed that the considered incoherent space-charge tune-shift is valid for a beam with infinitesimal coherent oscillations, as it is usually assumed in the perturbation theory of coherent instabilities. The study of the joint effect of octupoles and space charge for a beam with large coherent oscillations (decoherence effect) is more involved, as in this case the action variables for the incoherent space-charge tune shift are different from those with octupoles.

For an assumed particle density $n(x, y)$, Poisson’s equation provides the basis for obtaining the space-charge field components, assuming non relativistic beams and ignoring magnetic forces. It is well-known that Poisson’s equation can be integrated explicitly for the special case of ellipsoidal symmetry. Indeed, the solutions to Poisson’s equation for the electric field components from a distribution with elliptical symmetry, i.e. $n(x, y) = n(x^2 / x_m^2 + y^2 / y_m^2)$, assuming zero density outside the ellipse $x^2 / x_m^2 + y^2 / y_m^2 = 1$, are given by [42]

$$E_x = \frac{\varepsilon n_0}{2} \int \frac{x^2}{x_m^2 + y^2} \left( \frac{2}{y_m^2} \right) \left( \frac{2}{y_m^2} \right)^{-1/2} ds ,$$

$$E_y = \frac{\varepsilon n_0}{2} \int \frac{y^2}{x_m^2 + y^2} \left( \frac{2}{y_m^2} \right) \left( \frac{2}{y_m^2} \right)^{-1/2} ds .$$

Here, $\varepsilon_0$ is the permittivity of free space. Considering a round beam $(x_m = y_m)$ with the quasi-parabolic distribution function of Eq. (36) yields a particle density

$$n(x, y) = n_0 \left( 1 - \frac{x^2 + y^2}{x_m^2} \right)^3 .$$

Integrating $n(x, y)$ over the beam gives the total number of particles per unit length, which yields $n_0 = 2 N_b / (B \pi^2 R x_m^2)$, where $R$ is the average machine radius, $B = \sqrt{2 \pi} \sigma / (2 \pi R)$ is the bunching factor (considering a Gaussian longitudinal profile), with $\sigma_z$ the rms bunch length, and $x_m = \sqrt{10} \sigma_z$. The Lorentz force $F$, experienced by the particle located at the position $(x, y)$ is shown in Fig. 20, and is given by

$$F_x = \frac{\sigma_x}{\gamma^2} \frac{\varepsilon^2 n_0}{2 \varepsilon_0} \left[ x \left( \frac{x^2 + y^2}{x_m^2} \right) \left( \frac{x^2 + y^2}{x_m^2} \right)^{-1/2} \left( \frac{x^2 + y^2}{x_m^2} \right)^{-1/2} \left( 4 x_m^2 \right) \right] .$$

Figure 20: Defocusing space-charge force $F_x$ vs. $x / \sigma$ and $y / \sigma$ for a beam with the quasi-parabolic distribution of Eq. (36).

For an approximate solution, the nonlinear $x$- and $y$-dependence of the force is converted into an amplitude

237
dependence of the particle’s tune using the method of the harmonic balance, which is an averaging process over the incoherent betatron motions [38]. The self-consistent nonlinear space-charge tune shift is finally given by

$$\Delta Q_{\text{inc}}(j_x, j_y) = \Delta_0 + \frac{1}{K_1(c_1, q)} \left[ S_1 + 5 \varepsilon \Delta_x K_2(c_1, q) + 5 \varepsilon \Delta_y K_3(c_1, q) \right],$$

with

$$\Delta_0 = \frac{N_b r_p}{5 \pi B \beta y^2 e_{\text{norm}}},$$

where, $j_x = J_x / (5 \varepsilon)$, $j_y = J_y / (5 \varepsilon)$, $r_p$ is the classical proton radius, and $e_{\text{norm}} = \beta y e$ is the transverse rms normalized emittance.

Knowing the expression of the nonlinear space-charge tune shift, the dispersion relation of Eq. (35) can then be expressed. This equation has not been solved yet. As discussed in Ref. [43] and seen in Fig. 21 (in the case of the LHC at injection), a reasonable approximation of the space-charge tune shift is given by taking into account only the linear terms in the betatron action variables $J_{x,y}$ (adapting the coefficients!). In this case, it is written

$$\Delta Q_{\text{inc}}^\text{coh}(j_x, j_y) = \Delta_0 + \Delta_a J_x + \Delta_b J_y .$$

The dispersion relation of Eq. (35) can then be solved analytically and is expressed as

$$\Delta Q_{\text{coh}} = \Delta_0 + \frac{1}{K_1(c_1, q)} \left[ S_1 + 5 \varepsilon \Delta_x K_2(c_1, q) + 5 \varepsilon \Delta_y K_3(c_1, q) \right],$$

with

$$K_1(c_1, q) = -\frac{1}{6 c_1^2 (c_1 - 1)^2}$$

$$\times \left[ (c_1 + q)^3 \log(1 + q) - (c_1 + q)^3 \log(c_1 + q) \right] \right.$$

$$\times \left[ (c_1 - 1) \left[ \frac{c_1 + q + 2 c_1 q + (2 c_1 - 1) q^2}{(c_1 + q)^3} \log(q) - \log(1 + q) \right] \right]$$

$$+ \left[ (c_1 - 1) q^2 (3 c_1 + q + 2 c_1 q) \log(q) - \log(1 + q) \right]$$

$$K_2(c_1, q) = -\frac{1}{24 c_1^2 (c_1 - 1)^2}$$

$$\times \left[ (c_1 - 1) \left[ c_1 - 3 c_1^2 - 2 c_1 (c_1 + 2) q + c_1 (-11 + 5 c_1) q^2 \right] \right.$$

$$\times \left[ 2 \frac{1 + c_1 (-5 + 3 c_1)}{(c_1 + q)^3} \log(1 + q) \right]$$

$$+ 2 \left[ (c_1 + q)^3 \log(c_1 + q) \right]$$

$$\times \left[ (c_1 - 1) q^2 (4 c_1 + q + 3 c_1 q) \log(q) - \log(1 + q) \right]$$

$$K_3(c_1, q) = -\frac{1}{24 c_1^2 (c_1 - 1)^2}$$

$$\times \left[ (c_1 - 1) \left[ c_1^2 (1 + c_1) + 6 c_1^2 q + 3 c_1 (1 + c_1) q^2 \right] \right.$$

$$\times \left[ 2 \frac{1 + (1 + c_1) c_1}{(c_1 - 1 - c_1) c_1} \log(c_1 + q) \right]$$

$$+ 2 \left[ (c_1 + q)^3 (c_1 - q + 2 c_1 q) \log(1 + q) - 2 (c_1 + q)^3 \right]$$

$$\times \left[ (c_1 - q + 2 c_1 q) \log(c_1 + q) \right]$$

$$\times \left[ 2 (-1 + c_1) q^3 \log(c_1 + q) - \log(1 + q) \right]$$

$$\times \left[ (c_1 - q + 2 c_1 q) \log(1 + q) - 2 (c_1 + q)^3 \right]$$

$$\times \left[ (c_1 - q + 2 c_1 q) \log(c_1 + q) \right]$$

where $c_1 = b_1 / a_1$, $a_1 = a + \Delta_a$, $b_1 = b + \Delta_b$, $S_1 = -5 \varepsilon a_1$, and $q = (Q_c - Q_{\text{coh}} - m Q_{\text{inc}} - \Delta_0) / S_1$. 

Figure 21: 2D tune footprint for (upper left) the self-consistent and (upper right) the approximate ($\Delta_0 \approx -1.1 \times 10^{-3}, \Delta_a \approx 18127$ and $\Delta_b \approx 12948$) space-charge tune shift (the two plots are superimposed in the middle), and 3D tune footprint with the “shadows” on the horizontal and vertical tune axes for (lower left) the self-consistent and (lower right) the approximate tune shift. The position of the low-intensity small-amplitude working point ($Q_{\text{coh}} = 64.31$ and $Q_{\text{inc}} = 59.32$) corresponds to the upper right corner.
In the case of the LHC at injection (450 GeV/c), the nominal beam emittance is \( \varepsilon = 8.7 \text{ nm} \), and the maximum permitted octupole spread (compatible with an adequate dynamic aperture) yields the corresponding values of the anharmonicities \( a \approx \pm 7164 \) and \( b \approx \pm 4647 \). Making the numerical computation for the nominal LHC beam parameters gives \( \Delta_0 = -1.1 \times 10^{-5} \), \( \Delta_a \approx 18127 \) and \( \Delta_b = 1948 \).

Four stability diagrams are represented and compared in Fig. 22 with the nominal LHC beam parameters and maximum permitted octupolar strength: \( a > 0 \) corresponds to the case with octupoles only and positive horizontal detuning, \( a < 0 \) corresponds to the case with octupoles only and negative horizontal detuning, \( a > 0 + \text{ SC} \) corresponds to the case with both space-charge and positive horizontal detuning for the octupoles, \( a < 0 + \text{ SC} \) corresponds to the case with both space-charge and negative horizontal detuning for the octupoles. The evolution of these four stability diagrams with decreasing space-charge and octuplar strength is shown in Figs. 23 and 24 respectively. It is seen that when space-charge (i.e. intensity) decreases, the stability diagrams converge to the ones found by Berg and Ruggiero without space charge [35]. When the octupolar strength is reduced, the stability diagrams converge to each other and to zero, as predicted by Möhl and Schönauer [38].

Figure 22: Stability diagrams for the nominal LHC parameters at injection with maximum possible octupolar strength.

Figure 23: Evolution of the stability diagrams of Fig. 22 with decreasing space-charge (intensity): (a) \( N_s/2 \), (b) \( N_s/4 \), (c) \( N_s/10 \), (d) \( N_s/100 \).
The last missing important ingredient in this theory is the longitudinal variation of the transverse space-charge forces, which will fill the gap in the tune diagram (see Fig. 25, which has to be compared to Fig. 21 (upper left) obtained when the longitudinal variation of the transverse space-charge forces is neglected), and thus increase the stability region (on the right-hand side).

In conclusion, the above results should give a reasonable picture for beams with flat longitudinal profiles. In the usual case of parabolic or Gaussian bunches, the real stability region will be larger than predicted here. The following results are expected: the height of the stability diagram should be given by the octupoles (as observed before) and the width should be given by the small-amplitude incoherent space-charge tune shift for negative coherent tune shifts (as also observed before) and by the octupoles for positive coherent tune shifts.

Figure 24: Evolution of the stability diagrams of Fig. 22 with decreasing octupolar strength: (a) \( \Delta Q_{\text{oct,spread}}/2 \), (b) \( \Delta Q_{\text{oct,spread}}/4 \), (c) \( \Delta Q_{\text{oct,spread}}/10 \), (d) \( \Delta Q_{\text{oct,spread}}/50 \). Note the change of vertical scale for (c-d).

The stability of the longitudinal bunched-beam coherent mode \( m = \ldots, -1, 0, 1, \ldots \) can be discussed from the general dispersion relation [44]

\[
I_m(\omega) = \Delta \omega_{m,m}^2 ,
\]

(51)

where \( I_m(\omega) \) is the dispersion integral given by

\[
I_m(\omega) = \int_0^\infty \frac{\dot{\varphi}_0(\hat{t})}{\omega - \omega_0 - \dot{\varphi}_0(\hat{t})} \, d\hat{t}.
\]

(52)

Here, \( \Delta \omega_{m,m} \) is the coherent synchrotron frequency shift of Eq. (9), and \( g_0(\dot{t}) \) is the stationary distribution of the synchrotron oscillation amplitude \( \dot{t} \).

The stability diagram for the smooth distribution function \( g_0(\dot{t}) \propto (1 - \dot{t}^2)^3 \) used by Sacherer [44] is represented in Fig. 26, as well as the one corresponding to his “approximate” stability criterion \( S \geq 4 |\Delta \omega_{m,m}^2|/\sqrt{|m|} \) (following the example of Keil and Schnell for coating beams [27]). Sacherer approximated the stability boundaries by semi-circles). The case of a capacitive impedance below transition or inductive impedance above transition corresponds to \( \text{Re}(\Delta \omega_{m,m}^2/S) > 0 \) and \( \Delta \omega_{m,m}^2 < 0 \). Here \( S \) is the full spread between the centre and the edge of the bunch. As can be seen from Fig. 27, a good approximation of the frequency spread is given by [45]

\[
S = \left( 1 + \frac{5}{3} \tan^2 \phi_k \right) \frac{\pi^2}{16} (hB)^2 \omega_k .
\]

(53)
Consider the case of a capacitive impedance below transition or inductive impedance above transition. It is often said that the coherent synchrotron frequency remains the same as the unperturbed small-amplitude synchrotron frequency (coherent and incoherent effects subtract). As the incoherent frequency spread is moving downwards the following question is raised: how can the beam be stabilized by increasing the synchrotron frequency spread $S$, as it seems to be impossible, even for a very large frequency spread (see Fig. 28)?

An answer to this question was given by Besnier [46] for a parabolic distribution function. However this distribution introduces some pathologies in the stability diagram due to its sharp edge. As a consequence, no stable region was predicted in the case of a capacitive impedance above transition or inductive impedance below transition (see Fig. 29).

In the case of an “elliptical spectrum”

$$
\frac{d^2}{d\hat{\tau}^2} g_0(\hat{\tau}) \propto \sqrt{1-(2 \hat{\tau}^2-1)^2},
$$

the dispersion relation writes [47]

$$
S \geq \frac{4}{|m|} |\Delta \omega'_{\infty,m}|,
$$

Here, the coherent synchrotron frequency shift has been written $\Delta \omega'_{d1} = U - j V$. Motions $e^{j \omega t}$ are considered, which means that the beam is unstable when $V > 0$. Furthermore, the usual case where the resistive part of the impedance is small compared to the imaginary part, is assumed, i.e. $V \ll |U|$. Following Besnier’s approach, Fig. 29 is obtained.

The two plots of Fig. 29 are very similar, except that, contrary to Besnier, stability for the case of an inductive impedance below transition or a capacitive impedance above transition is also predicted with the elliptical spectrum. It is seen in Fig. 29 that in the absence of frequency spread ($S = 0$ and thus $k = 0$), the coherent synchrotron frequency $\omega_{c11}$ is close to the unperturbed
small-amplitude synchrotron frequency $\omega_{s0}$. When the synchrotron frequency spread increases, the coherent synchrotron frequency $\omega_{s1}$ moves closer and closer to the incoherent band (stable region). The two possible cases are represented in Fig. 29: the case of a capacitive impedance below transition or inductive impedance above transition corresponds to $U > 0$ and $\Delta \omega_s < 0$ (and thus $\omega_s < \omega_{s0}$), and the case of a capacitive impedance above transition or inductive impedance below transition corresponds to $U < 0$ and $\Delta \omega_s > 0$ (and thus $\omega_s > \omega_{s0}$). Beam stability is obtained when the coherent synchrotron frequency $\omega_{s1}$ enters into the incoherent band, i.e. when $\omega_{s1} = \omega_s$ for the case of a capacitive impedance below transition or inductive impedance above transition, and when $\omega_{s1} = \omega_s - S$ for the case of a capacitive impedance above transition or inductive impedance below transition. In both cases, the stability limit is reached for $k = 4$, i.e. $S = 4 |U |$, which is Sacherer's stability criterion (in the usual approximation $V \ll |U |$). Note that it is also the same stability criterion as the one used in Ref. [48] and derived in Ref. [49] (with the approximation $3/\pi \approx 1$).

Taking into account the PWD, the following stability criterion is obtained [47]

$$I_h \leq \left( 1 + \frac{5}{3} \tan^2 \phi_{s0} \right) \frac{3 \pi^2}{12} \times \frac{R_k}{\Omega_0} \times F_{\text{PWD}},$$

(55)

with

$$F_{\text{PWD}} = \sqrt{\frac{1}{2} \left( a + \sqrt{a^2 + 4} \right)},$$

(56)

$$a = \frac{9}{32} \left( 1 + \frac{5}{3} \tan^2 \phi_{s0} \right) h^2 B_k^2 Sgn \left( \cos \phi_{s0} \right) \frac{j}{Z_i(\frac{p}{\omega_{s0}})} \left| \frac{Z_i(\frac{p}{\omega_{s0}})}{p} \right|_{l=1}.$$

(57)

**Feedbacks**

An electronic feedback system is often used to damp coupled-bunch instabilities both in the longitudinal and transverse planes. Recently, it was found to help also for the head-tail instability in the Tevatron [50].

**Linear Coupling Between the Transverse Planes**

In the absence of both linear coupling and frequency spread, and below the mode-coupling threshold, the stability condition for the $m$th mode is $\text{Im} (\Delta \omega_{s,m}) \geq 0$, where $\text{Im} (\cdot)$ stands for imaginary part.

In the presence of linear coupling, but without frequency spread and below the mode-coupling threshold, the following necessary condition for stability of the $m$th mode is obtained [51]

$$V_m^m + V_m^m \leq 0,$$

(58)

where $V_m^m = -\text{Im} (\Delta \omega_{s,m})$ are the transverse instability growth rates. If Eq. (58) is true, it is possible to stabilize this mode by increasing the skew gradient and/or by working closer to the coupling resonance $Q_h - Q_v = l$. The stabilizing values of the modulus of the $l$th Fourier coefficient of the skew gradient are given by

$$\left| \bar{K}_h(l) \right| \geq \frac{2 \sqrt{|Q_h Q_v^m \Omega_0^m|}}{R_k^2 \Omega_0} \left[ \left( V_m^m + V_m^m \right)^2 + \frac{\Omega_0^m (Q_h - Q_v - l)^2}{4} \right]^{1/2},$$

(59)

where $Q_{h,v} = (\omega_{s0,0} + U_m^m)/\Omega_0$ are the horizontal and vertical coherent tunes in the presence of wake fields ($U_m^m = \text{Re} (\Delta \omega_{s,m})$), where $\text{Re} (\cdot)$ stands for real part, but in the absence of coupling. Notice that, when Eq. (58) is verified, it is verified for “any” intensity. The shape of the stable region is represented in Fig. 30. Furthermore, linear coupling has also a beneficial effect on the TMCI [52].

$$\left| \bar{K}_h(l) \right| \geq \frac{2 \sqrt{|Q_h Q_v^m \Omega_0^m|}}{R_k^2 \Omega_0} \left( V_m^m + V_m^m \right)\left( Q_h - Q_v - l \right)^2^{1/2},$$

Figure 30: Shape of the stable region in the presence of linear coupling when the sum of the transverse instability growth rates is negative.

A first experiment was performed with a CERN PS proton beam in 1997 (without Landau octupoles). The stability boundary was measured and compared to theoretical predictions (see Fig. 31) [9]. A second experiment was performed with a CERN PS proton beam in 1999 without Landau octupoles (see Fig. 5). As can be seen from Fig. 5, the sum of the transverse instability growth rates for the head-tail mode number $|m| = 6$ (in the absence of linear coupling) is negative, which means...
that this mode can be stabilized by linear coupling. This is indeed what was observed (see Fig. 32). Note that the PS beam for LHC is stabilized by linear coupling only (i.e. with neither Landau octupoles nor transverse feedback). The relation between the normalised skew gradient, as deduced from tune separation measurements, and the current in the skew quadrupoles is given in Fig. 33.

![Figure 31: Measured and theoretical stability boundaries for a CERN PS beam in 1997.](image)

![Figure 32: Intensity of the CERN PS ring (in units of 10^10 protons) vs. time (in ms). (a) Without linear coupling, i.e. I_{skew} = 0.33 A. (b) With a linear coupling corresponding to I_{skew} = -0.4 A, and a working point (Q_h = 6.22, Q_v = 6.25).](image)

In the presence of both octupoles and linear coupling between the transverse planes, the situation is more involved [8]. To clearly see the effect of linear coupling on the Landau damping mechanism of transverse coherent instabilities let's consider the case of a beam without frequency spread in the horizontal plane and without wake field in the vertical one (but assuming an elliptical vertical betatron frequency spread). In this case, the distribution function is defined by

$$\rho_y(\omega_y) = \frac{2}{\pi \Delta \omega_y} \sqrt{\Delta \omega_y^2 - (\omega_y - \omega_{y,0})^2},$$

where $\Delta \omega_y$ is the half width at the bottom of the distribution. A stability criterion can be derived and is given by

$$\left| Q_h - Q_v - l \right| \leq \frac{1}{\Omega_0}$$

$$\times \left( \frac{\Delta \omega_y^2}{4} \kappa^2 - \Delta \omega_y^2 \kappa + \Delta \omega_y^2 - V_{eq}^2 + \frac{4V_{eq}^2}{\kappa} - \frac{4V_{eq}^2}{\kappa^2} \right)^{1/2},$$

where

$$\kappa = \frac{|K_y(l)|^2 R^4 \Omega_0^4}{\Delta \omega_y \omega_{y,0} \omega_{y,0}}$$

The stability boundary relates the coherent tune separation to the linear coupling strength. If $\Delta \omega_y < V_{eq}$, then it is impossible to stabilize the beam by coupled Landau damping: there is not enough Landau damping which can be transferred to the unstable plane. The minimum frequency spread that can stabilize the beam is $\Delta \omega_y = V_{eq}$. In this case, there is only one condition for stability which is $Q_h - Q_v - l = 0$ and $\kappa = 2$. If $\Delta \omega_y > V_{eq}$, one can plot the curve describing the stability boundary. For example, in the case where $\Delta \omega_y = 2V_{eq}$, the absolute value of the tune separation $|Q_h - Q_v - l|$ vs. the normalised coupling strength $\kappa$ is represented in Fig. 34. It can be seen that below and above certain values of linear coupling strength,
stabilization is impossible whereas for intermediate values stabilization is possible even with some tune split \( Q_h - Q_v - l \). If the coupling is too small, there is not enough Landau damping transferred to the unstable plane. If the coupling is too large, the coherent frequencies fall outside the incoherent frequency spreads and Landau damping can’t exist. In the general case (with wake fields and frequency spreads in both transverse planes), an approximate stability criterion has been given [51]. The physical picture is shown in Fig. 35.

![Figure 34](image1.png)

**Figure 34**: Absolute value of the coherent tune separation \( |Q_h - Q_v - l| \) at the stability boundary vs. the normalised linear coupling strength \( \kappa \), for \( \Delta \omega_y = 2\nu_{eq} \).

![Figure 35](image2.png)

**Figure 35**: Transfer of frequency spread to Landau damp the horizontal plane. In the case where \( \nu_{eq} \ll \nu_{eq} \), which normally requires a large frequency spread for Landau damping in both planes, the e.g. vertical plane is stabilized by Landau damping (\( \nu_{eq} \) essentially) and the horizontal plane is stabilized by a judicious choice of linear coupling strength and tune distance from the resonance (the effect of \( \nu_{eq} \) is compensated).

As seen before, a too strong coupling can be detrimental since it may shift the coherent tunes outside the incoherent spectrum and thus prevent Landau damping. Considering the simplest case where the tune distributions and complex coherent tune shifts are the same in both transverse planes, the stability criterion (from an elliptical spectrum) is given by [53]

\[
\Delta \omega_{\text{spread}} \geq 2 \sqrt{\left( U_{\text{eq}} + \left| C \frac{\Omega_0}{2} \right|^2 \right)^2 + \nu_{\text{eq}}^2}.
\]  

Here, \( \Delta \omega_{\text{spread}} = \Omega_0 \Delta \nu_{\text{spread}} \) is the transverse betatron frequency spread (half width at the bottom of the elliptical spectrum), and \( |C| = \left| \Delta \nu_{\text{normal modes}} \right| \) is the modulus of the general complex coupling coefficient (it is the normal mode tune difference of the coupling resonance). If \( |C| = 0 \), i.e. in the absence of linear coupling, the one-dimensional (Keil-Zotter) stability criterion [54] is recovered. If \( |C| \) increases, the betatron frequency spread has to be increased according to Eq. (63). Therefore, any coupling is bad in that particular case since the stability condition of Eq. (63) is more restrictive than the one-dimensional stability criterion. When \( |C|\nu_{eq}/2 \gg \nu_{eq} \) and \( |C|\nu_{eq}/2 \gg \nu_{eq} \), Eq. (63) reduces to the following simple stability criterion

\[
\Delta \nu_{\text{spread}} \approx \left| \Delta \nu_{\text{normal modes}} \right|.
\]  

To prevent the instability from developing, the normal mode tune difference on the coupling resonance has to be kept smaller than the half width at the bottom of the tune distribution.

Linear coupling should also be used with great care when the beam is not round, as it will lead to a sharing [55], or exchange [56], of the transverse emittances (therefore modifying the stability diagram for Landau damping) and possibly beam loss due to vertical acceptance limitation. Two regimes are indeed possible in the presence of linear coupling: the static one (when the beam is injected directly into the stop-band, and the tunes are kept constant) and the dynamic one (when the beam is injected outside the stop-band and then the working point is shifted across the resonance). In the static case, the transverse emittances are given by [55]

\[
\varepsilon_x = \varepsilon_{x,0} - \left( \varepsilon_{x,0} - \varepsilon_{y,0} \right) \frac{|C|^2 / 2}{\Delta^2 + |C|^2},
\]

\[
\varepsilon_y = \varepsilon_{y,0} + \left( \varepsilon_{x,0} - \varepsilon_{y,0} \right) \frac{|C|^2 / 2}{\Delta^2 + |C|^2},
\]

where \( \varepsilon_{x,y,0} \) are the initial uncoupled horizontal and vertical emittances, and \( \Delta = Q_h + l - Q_v \). It can be seen
from Eqs. (65) and (66) that in the presence of coupling, there is a sharing of the emittances as coupling increases.

Figure 36: Measured normal mode tunes vs. time.

Figure 37: Transverse physical emittances vs. time for $A_{1.0} \approx \text{skewI}$ (upper), $A_{5.0} \approx \text{skewI}$ (centre), and $A_{5.1} \approx \text{skewI}$ (lower). The dashed lines denote the theoretical values and the solid lines the measured ones.

As can be seen from Eqs. (67) and (68), in addition to the emittance sharing, an emittance exchange is also predicted after the resonance crossing. These formulae have been confirmed by simulations [58]. Equations (67) and (68) have been compared to measurements in 2002 in the CERN PS [59], which were performed by programming the transverse tunes to slowly exchange their values within 100 ms on the injection flat-bottom. The measured normal mode tunes are shown in Fig. 36 for a particular coupling strength, where the closest-tune-approach is given by $|C|$. During the 100 ms needed to exchange the transverse tunes, the transverse emittances were measured several times with a Wire Scanner, and for different coupling strengths (see Fig. 37). As can be seen from Fig. 37, the horizontal and vertical emittances are shared until full coupling (on the coupling resonance), where the emittances become equal. Then the emittances are exchanged in good agreement with the theoretical predictions. The remaining small difference between the measurements and the theoretical predictions (in particular in the first case with very small linear coupling strength) may be due to the so-called Montague resonance [60], which is an intrinsic space-charge driven fourth-order resonance $022 = -v_h Q h Q Q [61]$.

REFERENCES


INTENSITY LIMITATIONS BY COMBINED AND/OR (UN)CONVENTIONAL IMPEDANCE SOURCES

G. Rumolo∗
GSI, Darmstadt, Germany and Università di Napoli “Federico II”, Napoli, Italy

Abstract

The main intensity limitations in future hadron or positron rings are expected to come from unconventional loss mechanisms. First, we will present a review of the single bunch collective effects driven by electron cloud. The wake field generated by an electron cloud can lead to transverse beam instability and loss over a few turns. The electron cloud wake field is highly unconventional, since the shape of the wake depends on the cloud distribution, and therefore on which part of a passing bunch is generating it. This feature leads to the definition of a generalized impedance, which has a double frequency dependence and can be used in TMCI calculations for the accurate estimation of the instability threshold. Space charge and conventional broad band impedance can also interplay with the electron cloud to trigger stronger instabilities.

The second part of the paper will address the features of collective head-tail motion in a barrier bucket under the action of a transverse broad band impedance. Similarities and differences with regular and strong head tail instability for a conventional bunch shape will be highlighted.

INTRODUCTION

The electromagnetic interaction of a high intensity beam with the surrounding environment inside an accelerator ring is recognized to be responsible for unstable collective motion and unwanted beam loss. When the beam intensity is sufficiently high the electromagnetic field self-generated by the beam perturbs the external prescribed fields and acts back on the beam, perturbing in turn its motion. Under unfavourable conditions, the perturbation on the beam further enhances the perturbation on the fields, and an unstable mechanism is initiated. The subject of collective instabilities in accelerators has been studied since the early 1960s [1, 2]. The impact of the understanding of collective instability mechanisms in determining the ultimate performance of an accelerator defines the importance of the subject. Each accelerator, when pushed for performance, encounters some intensity limit, which needs to be understood and cured before moving on to the next limit.

The concepts of wake fields and impedance [3, 4] have been introduced and are used to describe this class of phenomena. Every kind of interaction of the beam with itself and with the environment as defined by its geometrical and physical properties, is what we classify in this paper as “conventional impedance sources”. They include: space charge, resistive wall (including the effects of small holes, surface roughness, finite conductivity and finite thickness), narrow- and broad-band impedance (modeling the whole accelerator or single objects like cavities, pick up electrodes, kickers, vacuum ports, pipe discontinuities). All these impedance sources have been extensively studied over the years, and several methods for impedance calculation (Maxwell solvers in complex structures) and measurements (wire, beam-based) have been developed [5]. References [6, 7] are very comprehensive overviews on analytical calculation of conventional impedances and wake fields. Recent work has been or is still currently being carried out on the generalization of space charge formulae [8] and resistive wall in non-ideal conditions (flat chamber, non-relativistic beam, holes, finite thickness, multi-layer chamber) [9, 10, 11, 12, 13, 14].

All the types of electromagnetic interaction which do not belong to the class previously defined, can be identified as “unconventional impedance sources”. They include: interaction of the beam with a medium (e.g., non-neutral one component electron plasma, magnetized or not, such as an electron cooler [15, 16] or an electron cloud [17, 18]) and with the electromagnetic field radiated by the beam itself in the arcs (Coherent Synchrotron Radiation) [19]. The wake fields associated to these interactions have usually significantly different, and more complex, features than conventional wake fields. For instance, an electron cloud as well as a detuned electron cooler generate an amplitude and location dependent wake field (see the following). Since particles move with velocity lower than c on a curved trajectory, the synchrotron radiation that they emit causes a “forward wake field”, which affects witness particles situated in front of the source particle.

Even if since the beginning of the 1970s the question of two stream instabilities driven by electrons from residual gas ionization trapped around coasting hadron beams had been investigated, it was not before 1999 that the concept of “two-stream transverse impedance” was introduced [20]. Nevertheless, the main interest has moved meanwhile to the problems caused by multipacting of electrons in machines operating with trains of short positron or hadron bunches. The detailed study of the trailing field caused by the passage of a bunch through an electron cloud highlights all its unconventional features (Section 2.1). A more general definition of impedance is required to fit this case into the impedance formalism (Section 2.2).

Transverse coherent motion of a bunch having a uniform profile in a barrier bucket is the subject of the second part of this paper. Dipole and envelope oscillation modes are studied in detail as functions of the bunch current and a com-

∗G.Rumolo@gsi.de
parison with a regular Gaussian bunch inside a sinusoidal bucket is shown (Section 3.1). Strong and regular head-tail instabilities in barrier buckets are then investigated in Section 3.2, and conclusions as well as suggestions for possible future work are drawn in Section 4.

![Figure 1: Electron cloud wake field of an SPS bunch: dependence on the location of the displaced “source slice”. The top picture shows the averaged wake functions, the bottom picture shows the wake functions on axis.](image1)

**ELECTRON CLOUD WAKE FIELDS**

*Parameter dependence of the electron cloud wake fields*

Electron cloud wake fields are usually found by macro-particle simulation. One bunch slice is displaced (for instance, vertically by a small amount \( \Delta y \)), and the electron cloud response is evaluated in terms of electron field on axis (\( x = y = 0 \)). Normalizing this field by the amount of displacement and the number of particles contained in the displaced slice, we obtain the dipole wake function on axis (in \( \Omega s^{-1} m^{-1} \), after multiplication by the factor \( m\gamma c^2/e^2 \)). As the field on axis is not directly related to the force exerted by the cloud on slices that follow the displaced one, we can in addition evaluate an averaged dipole wake function from the net force caused by a displaced slice on later portions of the beam. In this case, instead of looking only at the field on axis, we calculate the overall force exerted by the distorted cloud on the particles contained in one slice, and then divide by the total charge in the slice to obtain an effective electric field. Shapes in the two cases appear quite different, as shown in Figs. 1. Note that the two definitions of the wake field would lead to the same result for a conventional wakefield. Wake functions on axis reach much larger values and exhibit a spiky structure that is smoothed out to a quasi-sinusoidal profile when the integration over the bunch slice is carried out. These plots correspond to an almost round beam in an SPS field-free region and are calculated for a longitudinally uniform bunch distribution. It is worth noting two remarkable features that make the electron cloud wake field differ from a conventional dipole wake field: first, the two definitions of the wake field given above do not lead to the same result, and second, the shape of the wake depends on the longitudinal location of the displaced slice.

![Figure 2: Electron cloud averaged wake field of an SPS bunch: dependence on the electron cloud distribution. The top picture shows the horizontal component, the bottom picture shows the vertical component.](image2)

In a region with a dipolar magnetic field (bending magnet) electrons accumulate in the form of one or two vertical stripes (generally depending on the bunch current). The horizontal wake tends to vanish and the vertical one becomes also weaker. Figures 2 show the averaged dipole wake functions (\( x \) and \( y \)) of a Gaussian bunch distribution in a bending region for different initial electron distributions (one stripe \( 4\sigma_x \) wide, and then two \( 2\sigma_x \) wide stripes...
located at increasing distances from the axis). The frequency of the wake decreases as the separation between the two electron stripes is increased. The horizontal wake is almost independent on the distance of the stripes from the beam axis, whereas the vertical gets weakened with increasing distance.

The electron cloud wake field also strongly depends on the transverse distribution of the protons. Figure 3 shows that for a uniform transverse proton distribution the wake function (both the averaged one and the one on axis) does not get damped far from the source slice. In terms of broadband impedance model, this would correspond to considering a resonator with an infinite merit factor, $Q$. Non-uniformity of the bunch transverse distribution leads to frequency spread of the electron oscillation and results in decoherence of the cloud response. For a Gaussian transverse distribution, the memory of the electron cloud only extends over a finite length and a broad-band resonator model with finite $Q$ is therefore possible [17].

The dependence of the electron cloud wake field on the boundary conditions has also been studied. For an SPS bunch and uniform transverse distribution of the protons, we can deduce from Figs. 4 that whether we account or not of perfectly conducting boundary conditions on a rectangular pipe does not make a large difference in the resulting wake field. This is true as long as the bunch transverse sizes are much smaller than the pipe size and the beam is on axis. As an unstable motion is initiated, which makes the beam centroid oscillate to large amplitudes and/or the beam transverse size grow closer to the pipe size, the difference becomes significant and the correct boundary conditions have to be considered in electron cloud simulation studies.

Figure 3: Electron cloud wake field of an SPS bunch: dependence on the proton transverse distribution. The top picture shows the averaged wake functions, the bottom picture shows the wake functions on axis.

The dependence of the electron cloud wake field on the boundary conditions has also been studied. For an SPS bunch and uniform transverse distribution of the protons, we can deduce from Figs. 4 that whether we account or not of perfectly conducting boundary conditions on a rectangular pipe does not make a large difference in the resulting wake field. This is true as long as the bunch transverse sizes are much smaller than the pipe size and the beam is on axis. As an unstable motion is initiated, which makes the beam centroid oscillate to large amplitudes and/or the beam transverse size grow closer to the pipe size, the difference becomes significant and the correct boundary conditions have to be considered in electron cloud simulation studies.

Generalized impedance for electron cloud wake fields

When defining wake functions in beam physics, one normally assumes that the effect of a source particle on a witness particle only depends on the distance between the two particles, $W(z - z')$. For more general situations, e. g. for the electron cloud response to dipole perturbations, translation invariance does not hold, as shown in the previous subsection. In this case, the wake function cannot be reduced to the form $W(z - z')$, but a generalized definition must be given, $W(z, z')$ with $z > z'$. First, we introduce then the generalized impedance as the
full Fourier transform of this wake function,

\[ Z(\omega, \omega') = 2\pi \int \int W(z, z') \exp \left[ i(\omega - \omega') z_t \right] dz dz', \]

where the particular case \( Z(\omega, \omega') = 2\pi \delta(\omega - \omega') Z(\omega) \) is recovered from the conventional wake \( W(z - z') \).

Following the kinetic approach through the Vlasov equation, the integral equation that gives the frequency shifts of the different modes has the form [21]:

\[ \left( \frac{\Omega - \omega}{\omega} - i \right) a_k = -\frac{N_0 T_{rc}}{2 \gamma \omega \omega_s T_0} \sum_{l',k'} M_{lk,l'k'} a_{l'k'}, \]

where the mode coupling matrix \( M_{lk,l'k'} \) is expressed for a Gaussian bunch as

\[ M_{lk,l'k'} = -\frac{N_0 T_{rc}}{\gamma \omega \omega_s} \frac{|l-l'| \epsilon(l) \epsilon(l')}{k! (|l| + k)! (|l'| + k')!} \]

\[ \times \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} \omega d\omega' Z(\omega + \omega_{l}, \omega + \omega_{l'}) \]

\[ \times \left( \frac{\omega_{l'}}{\sqrt{2c}} \right)^{|l|+2k} \left( \frac{\omega_{l} \omega_{l'}}{\sqrt{2c}} \right)^{|l'|+2k'} \exp \left[ -\left( \omega^2 + \omega'^2 \right) / 2c^2 \right]. \]

Using Eq. 2 together with Eq. 3, we can calculate the dependence of the frequency shifts of the different modes on beam intensity, and find existence and value of the threshold above which mode coupling, and therefore instability, occurs. For this purpose we need to know the generalized impedance.

An example of calculation of generalized impedance to be used in this model is given in the following. The wake function generated by displacing several slices of a Gaussian SPS bunch going through an electron cloud is evaluated numerically and shown in Fig. 5.

![Figure 5: Electron cloud average wake field of an SPS bunch calculated by displacing a “source slice” at several locations along the bunch.](image)

Figure 5: Electron cloud average wake field of an SPS bunch calculated by displacing a “source slice” at several locations along the bunch.

The calculation of the generalized impedance follows then in a quite straightforward way. If \( N \) is the number of slices in which the bunch was subdivided, the wake fields are stored in an \( N \times N \) matrix, on which subsequently a double Fourier transform is applied. Two example results are shown in Fig. 6 for the case of the CERN-SPS in Fig. 5 and for the GSI-SIS, where the electron cloud is given by the detuned electron cooler [22].

![Figure 6: Two frequency impedance \( Z(\omega, \omega') \) for the CERN-SPS (top) and for the GSI-SIS.](image)

**COLLECTIVE MOTION IN BARRIER BUCKETS**

The interest in studying the beam dynamics in barrier buckets has recently become stronger thanks to the progress in using induction cavities in accelerators to accelerate and confine particle bunches in square buckets [23]. Square bunch profiles can also be created in an accelerator ring by tuning several rf-systems on different multiples of the harmonic that is meant to be generated. In particular, the use of higher harmonics has been suggested for the LHC upgrade to create flat bunches which are about 10 times longer than the bunches in nominal operation. These bunches are not quite as long and as performant as the previously proposed “super-bunches” [24], but they are certainly able to deliver higher luminosity than in nominal operation [25]. The use of double harmonic rf-systems to flatten bunch profiles against space charge has been also envisaged at GSI in order to optimize the performances of the new accelerator rings SIS100 and 300 [26]. Even if a two harmonic rf-system is not a barrier bucket, the flattened force profile at the center of the bucket makes the longitudinal beam dynamics of particles in it closer to that of particles in a barrier bucket than in a conventional sinusoidal bucket, if the bunch is sufficiently short. Previous work on
headtail motion in a barrier bucket has been presented [27], in which a better stability of long uniform bunches against regular head-tail instability was highlighted.

**Dipole and envelope mode shifts**

In our study, we model the bunch particles as confined between two infinitely high voltage barriers (elastic reflection occurring at the bucket ends in zero time). The bunch feels then the action of a broad band impedance. It is subdivided into \( N \) slices and each slice feels the sum of the wakes of all preceding slices. The fields considered are dipole as well as quadrupole wakes for flat chambers (weighted with the Yokoya coefficients). The resulting centroid and envelope oscillations are calculated with the HEADTAIL code [28] and the motion is Fourier analysed in order to find out dipole and envelope modes, as well as mode dependence on the bunch intensity.

We have chosen to do our simulations using a relatively long flat SPS bunch (not really attainable in the machine), in order to compare the results with those known from the real Gaussian SPS bunch having the same longitudinal emittance. The parameters we have used in our simulations are summarized in Table 1.

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Symbol</th>
<th>Number</th>
</tr>
</thead>
<tbody>
<tr>
<td>Protons/bunch</td>
<td>( N_b )</td>
<td>0.2 to 2 ( \times 10^{11} )</td>
</tr>
<tr>
<td>Energy</td>
<td>( E )</td>
<td>26 GeV</td>
</tr>
<tr>
<td>Bunch length</td>
<td>( \sigma_z )</td>
<td>2.4 m</td>
</tr>
<tr>
<td>Mom. spread</td>
<td>( \delta )</td>
<td>2 ( \times 10^{-3} )</td>
</tr>
<tr>
<td>Mom. compaction</td>
<td>( \alpha )</td>
<td>1.92 ( \times 10^{-3} )</td>
</tr>
<tr>
<td>Transv. sizes</td>
<td>( \sigma_{x,y} )</td>
<td>1.2 mm</td>
</tr>
<tr>
<td>Tunes</td>
<td>( Q_{x,y} )</td>
<td>26.185/26.13</td>
</tr>
<tr>
<td>Chromaticities</td>
<td>( \xi_{x,y} )</td>
<td>0 to -1</td>
</tr>
<tr>
<td>Shunt imp.</td>
<td>( R_s )</td>
<td>20.40 M( \Omega )/m</td>
</tr>
<tr>
<td>Res. frequency</td>
<td>( \omega_r )</td>
<td>1.3 GHz</td>
</tr>
<tr>
<td>Qual. factor</td>
<td>( Q )</td>
<td>1</td>
</tr>
</tbody>
</table>

As the bunch intensity is scanned through the values reported in the Table above (from 2 \( \times 10^{10} \) to 2 \( \times 10^{11} \)), we can see how the tune line, obtained from the Fourier spectrum of the centroid motion of the bunch after a transverse kick, moves leftward (negative tune shift, as for conventional bunches). If we assume a round chamber, there is no difference in behaviour between the two transverse planes (Figs. 7).

Analyzing the envelope motion of the same SPS bunch inside a round chamber, one finds out that two modes are clearly visible in both planes. The first mode corresponds to twice the non-shifted tune and does not change its position with increasing intensity. The second line shows instead a growing amplitude with intensity, and also its position shifts toward smaller values as the bunch current increases. As the separation of the modes gets larger with larger currents, envelope instabilities of the transverse mode coupling type do not appear.

![Figure 7: Mountain ranges of horizontal and vertical tune shifts with increasing current (20 M\( \Omega \)/m): round chamber.](image7.png)

![Figure 8: Mountain ranges of horizontal and vertical envelope modes with increasing current (20 M\( \Omega \)/m): round chamber.](image8.png)
Fourier analysis of the centroid motion of the SPS bunch inside a flat chamber shows the typical features of the coherent shifts in flat chambers: the horizontal tune shift disappears, whereas the vertical tune shift stays practically unchanged. Figures 9 display the mountain range plots of the horizontal and vertical tune shifts as functions of intensity inside a flat chamber.

Due to the quadrupole wake field present in flat chambers because of the impedance, the envelope modes of a bunch in a flat chamber show evident differences with respect to the case of round chamber (Figs. 10). The horizontal envelope spectrum exhibits one single line, which has growing amplitude with current, as well as its position shifts towards larger frequencies when the bunch current is made to increase. The vertical spectrum still has two lines with amplitude increasing with current. They both shift towards smaller frequency values as the bunch current grows larger. The quicker shift of the line in lower frequency makes sure also in this case that no mode coupling occurs at the higher currents, since the separation between the two modes grows with the bunch intensity.

The linear tune shifts with intensity and parameter dependence are summarized in the Figs. 11 through 16. Figures 11 and 12 show that the slope of the tune shift line with bunch current is directly proportional to the shunt impedance of the broad band resonator and inversely proportional to the bunch length. Figures 13 prove that a flat chamber causes complete cancellation of the horizontal tune shift, whereas the vertical tune shift stays basically unchanged. The comparison with a Gaussian bunch in a sinusoidal bucket reveals interesting features (Fig. 14). At the low intensities, the dominant dipole mode of a conventional bunch corresponds to the tune line of a barrier bucket bunch and the shift with intensity occurs with the same slope (theoretical line). At the higher currents, other synchrotron modes become dominant for a conventional bunch, and the peak detection algorithm finds the tune by about one synchrotron frequency shift in the current range \(8 - 12 \times 10^{10}\), by about two synchrotron frequencies shift in the range \(12 - 14 \times 10^{10}\), and finally the strong head-tail instability sets in on a mode merging line for current values higher than \(15 \times 10^{10}\).

Envelope line positions versus bunch intensity are plotted in Figs. 15 and 16. Figures 15 show both main and secondary lines in \(x\) and \(y\) for an SPS flat bunch in a round chamber. Independently of the shunt impedance value of the broad band resonator used in our simulations, the main line does not move changing the bunch intensity. The slope of the secondary line is on the other hand proportional to the shunt impedance value, which is seen because the variation of the line position with intensity doubles when the shunt impedance gets doubled. In a flat chamber, the main line in the horizontal spectrum shifts toward higher frequencies when the current increases, and the slope is proportional to the shunt impedance, first of Figs. 16. The main line in the vertical spectrum shifts toward lower frequencies when the current increases, and again the slope is proportional to the shunt impedance, second of Figs. 16.
Figure 11: Horizontal and vertical tune shifts versus current. Dependence on the impedance.

Figure 12: Horizontal and vertical tune shifts versus current. Dependence on bunch length.

Figure 13: Horizontal and vertical tune shifts versus current. Dependence on chamber shape.

Figure 14: Horizontal and vertical tune shifts versus current. Barrier versus sinusoidal bucket.
Finally, the secondary line, which in this case only appears in the vertical spectrum, has a shift to the lower frequencies with the same slope as the secondary line from the envelope spectrum of a bunch in a round chamber. The slope is also in this case proportional to the shunt impedance value, as shown in the third of Figs. 16.

Head-tail instability

In conventional buckets, the transverse head-tail instability can appear:

- At whatever bunch current, when chromaticity is positive and the beam is below transition or chromaticity is negative and the beam is above transition. Typical rise times of this instability are in the order of a few synchrotron periods because the head-tail exchange through synchrotron motion is an essential ingredient in the excitation of the instability.
- Above a certain threshold current, independently of chromaticity. This instability, known as strong head-tail or transverse mode coupling instability, is of the beam break-up type (similar to what happens in linacs) and has typical growth times much shorter than the synchrotron period, because it must set in before synchrotron motion can cause enough phase mixing within the bunch so as to prevent it. The Vlasov formalism shows that the threshold current at which this instability appears, is determined by mode coupling.
geometries (round and flat). Rise times double as the shunt impedance of the resonator is halved or the bunch length is doubled. In a flat chamber, the rise times observed in the horizontal plane are twice the rise times in the vertical plane. Comparison with a Gaussian bunch in a sinusoidal bucket shows that the difference is in this case very tiny. The behaviour of a bunch in a barrier bucket with respect to head-tail instability in the weak regime does not differ much from that of a bunch in a sinusoidal bucket if the two bunches have the same longitudinal emittance.

Figure 17: Growth times of head-tail instability: dependence on impedance, chamber shape and bunch length.

The strong head-tail instability regime is not observed in barrier buckets. This should not be surprising, because modes are not distinguishable in a barrier bucket (the synchrotron side bands merge into an incoherent spread of the main tune line). Increasing the bunch current, we observe that the bunch enters a regime of slow growth, but neither a violent instability nor a clear threshold are identifiable. Figures 18 show the difference between a sinusoidal and a barrier bucket. A Gaussian bunch in a sinusoidal bucket clearly crosses a threshold, above which the centroid motion of the bunch becomes unstable and increases exponentially, whereas a flat bunch in a barrier bucket exhibits a slow growth with increasing intensity but no threshold exists.

Figure 18: Strong head-tail instability: in a Gaussian bunch in sinusoidal bucket there is a sharp threshold (bottom plot), in a barrier bucket there is a regime of slow growth (top plot).

CONCLUSIONS

Electron cloud wake fields and relative impedance, as well as coherent motion of a barrier bucket under the effect of a broad band impedance have been the subject of this paper. For an electron cloud generated by a bunched beam, the concepts of wake field and impedance require an extension of the conventional definitions due to the unconventional features summarized here below:

- The dipole wakefield of an electron cloud depends on the transverse coordinates \((x, y)\)
- Differently located displacements along a bunch create differently shaped wake fields
- The wake field depends:
  - Strongly:
    * On the initial electron distribution
    * On the bunch particle transverse distribution
  - Weakly:
    * On the boundary conditions for a wide pipe
    * On the electron space charge for low degrees of neutralization
- A description in terms of double frequency impedance $Z(\omega, \omega')$ is necessary for a correct Transverse Mode Coupling (TMC) analysis

Numerical tools to handle the calculation of $Z(\omega, \omega')$ have been developed, but both the dependence of the wake on the longitudinal shape of the bunch and the electron cloud wake field for long bunches (which might strongly depend on the trailing edge electron production and multiplication) are yet to be investigated.

An accurate analysis of dipole and envelope modes of a bunch in a barrier bucket leads to the following conclusions:

- The coherent tune shift $\Delta Q$ of a bunch in a barrier bucket as a function of the bunch current depends on
  - Shunt impedance (proportional)
  - Bunch length (inversely proportional) and momentum spread (proportional)
  - Chamber shape (only in $x$)
- The $\Delta Q$ in low current follows that of a usual bunch in a sinusoidal bucket with the same longitudinal emittance (theoretical line).
- Coherent envelope modes depend on the chamber shape:
  - Round chamber has two modes both in $x$ and $y$, one current dependent and one current independent.
  - Flat chamber has one mode in $x$ with a positive shift with increasing current, and two modes in $y$, both with a negative shift with current.

Concerning instabilities:

- High current: The threshold for strong head-tail instability is not found for bunches in a barrier bucket, but there is rather a regime of slow growth at high currents.
- Current independent: Regular head-tail instability driven by negative $Q'$ (above transition) exhibits similar features as for bunches in sinusoidal buckets.
  - Growth rates are proportional to the shunt impedance
  - The quickest instability occurs when $\omega_\xi = \omega_r$.
- In a flat chamber growth times in the $x$ direction are about double of the growth times in the $y$ direction
- Longer bunches slow down the instability

An analytical model (maybe few particles model or kinetic model based on Vlasov equation) is needed.

**ACKNOWLEDGMENTS**

The author would like to thank Oliver Boine-Frankenheim, Elena Shaposhnikova and Frank Zimmermann for interesting and fruitful discussion.

**REFERENCES**


[27] Y. Shimosaki, T. Toyama, K. Takayama, “Head-tail instability of a Super-bunch”, in Proc. of the ICFA Workshop on High Brightness and High Intensity Beams, Ref. [23]

Code comparisons and benchmarking with different SEY models in electron cloud build-up simulations

G Bellodi*, ASTeC Intense Beams Group, RAL, Chilton, Didcot OX11 0QX, UK

Abstract

Several phenomenological fits modelling the secondary electron emission process (yield and energy spectrum) are currently used in computer codes simulating electron cloud formation. Here we present a comparison of preliminary simulation results for two of these existing fits [1, 2] and a set of recent experimental measurements [3], when using identical or nearly identical input parameters.

INTRODUCTION

A number of computer codes has been developed in laboratories worldwide to simulate electron cloud formation and dissipation in accelerators. The codes’ underlying assumptions and physical models are often different and based on alternative experimental inputs. As part of an inter-laboratory code-code and code-experiment benchmarking effort re-launched after the ECLOUD’04 workshop [4], the present study explores the effect on simulation results of a reference case of using different descriptions of secondary emission properties for electrons impinging on the vacuum chamber walls. The two models currently used in the codes ECLoud [5, 6] and POSINST [7, 8] (and described respectively in [1] and [2]) have been compared, along with some tentative phenomenological fits of recent laboratory measurements obtained in surface science experiments realised at CERN [3].

MODEL DESCRIPTION

The two main quantities describing electron secondary emission are the yield per incident particle and the secondary energy spectrum. The yield $\delta$ is a function of the kinetic energy of the incident electron $E_0$, its incident angle $\theta_0$ and the type and condition of the surface material. $\delta$ is conventionally considered as consisting of two to three components, namely an elastic part (backscattered electrons), a true secondary and a rediffused part (electrons scattered in the material and reflected back):

$$\delta(E_0, \theta_0) = \delta_{true}(E_0, \theta_0) + \delta_{el}(E_0, \theta_0) + \delta_{red}(E_0, \theta_0).$$  \hfill (1)

The last component is not always included in all models though, owing to what some surface scientists consider too high a level of arbitrariness in the distinction from the true secondaries at low energies.

An analytical formula given in [7] has been adopted by both the ECLoud and POSINST codes to describe the true secondary component:

$$\delta_{true}(E_0, \theta_0) = \delta_{max} D(E_0/E_{max}),$$  \hfill (2)

$$D(x) = \frac{s(x)}{s - 1 + (x)^s},$$

where $s \approx 1.54$ for Copper and $\delta_{max}$ and $E_{max}$ include a dependency on $\cos \theta_0$. The expression for $D(x)$ was chosen as the simplest function satisfying the conditions: $D(1)=1$ and $D'(1)=0$, and that allowed a good fit to experimental data [9].

In the original version of ECLoud, the elastic component is then defined as a function $f$ of the true secondary [1]:

$$\delta_{el} = \frac{f}{1-f} \delta_{true},$$  \hfill (3)

where $\ln(f) = A_0 + A_1 \ln(E_0 + E_c) + A_2(\ln(E_0 + E_c)^2 + A_3(\ln(E_0 + E_c)^3$, with $A_0 \approx 20.7, A_1 \approx 7.08, A_2 \approx 0.48, A_3 \approx 0$ and $E_c \approx 56.9$ eV for incident energies $E_0$ below 300 eV. Fig. 1 shows the secondary emission yield curves that can be extracted from this model, when assuming $\delta_{max}=2.03$ and $E_{max}=262$ eV for the case of Copper.

In POSINST, on the other hand, the elastic component is modelled by:

$$\delta_{el} = P_{1,c}(\infty) + \left[\hat{P}_{1,c} - P_{1,c}(\infty)\right] e^{-(E_0 - E_c)/W}p, \hfill (4)$$

where $P_{1,c}(\infty) \approx 0.02$, $\hat{P}_{1,c} \approx 0.476$, $W \approx 60.86$ eV, $E_c=0$ eV and $p=1$. A rediffused component is also included, and is described by:

$$\delta_{red} = P_{1,r}(\infty) \left[1 - e^{-(E_0/E_r)^2}\right],$$  \hfill (5)

* g.bellodi@rl.ac.uk
where $P_{1,r}(\infty) \simeq 0.19$, $E_\nu \simeq 0.041$ eV and $r \simeq 0.104$. Fig. 2 shows these curves as a function of the primary electron energy for the case of perpendicular incidence.

Another difference between the codes is in the modelling of the secondary energy distribution: in ECLOUD the true secondary energy spectrum is described by a formula first presented in [10]:

$$dN_s/dE_s \propto \exp \left[ -0.5 \left( \ln (E_s/E_0) \right)^2 / \tau^2 \right],$$  \hspace{1cm} (6)

where $E_0$ and $\tau$ are fitting parameters assumed to be $\simeq 1.8$ eV and $\simeq 1$ respectively. Alternative descriptions might employ a Gaussian or Lorentzian distribution, falling to zero at $E_s=0$ eV. Reflected secondary electrons are assigned the same energy and angular distribution as the incoming electron, whereas the true secondaries are individually assigned an energy from a uniform distribution between 0 and 30 eV.

POSINST adopts a more complicated, mathematically self-consistent, probabilistic model, where every event is assigned a set of probabilities for the generation of electrons, following a Monte Carlo procedure where phenomenological fits are used. For the case $E_0 \gg E$, the true secondary energy spectrum is modelled by:

$$\frac{d\delta_s}{dE} = \sum_{n=1}^{\infty} n P_{n,t_s}(E_0) \frac{\left( E/\epsilon_n \right)^{n-1} e^{-E/\epsilon_n}}{\epsilon_n \Gamma(p_n)},$$  \hspace{1cm} (7)

where $P_n$ is the probability for the event and $p_n$, $\epsilon_n$ are fitting parameters. Reflected secondary electrons are again assigned the same energy (minor a small Gaussian smearing) and direction as the incident electron, while the rediffused electrons’ energy follows: $E = E_0 r^{0.7}$, where $r$ is a uniformly distributed number between 0 and 1.

**BUILD-UP SIMULATIONS FOR A REFERENCE CASE**

Electron cloud build-up simulations have been carried out using the two secondary emission models described above, both inserted in the ECLOUD code, for a machine reference case whose parameters, resembling those of the LHC proton beam, are listed in Table 1.

![Buildup results for a) the ECLOUD SEY model (in blue) or b) the POSINST-like SEY model in full (in red) or with some components suppressed (magenta and green).](image)

**Table 1: Input parameters for the study reference case.**

<table>
<thead>
<tr>
<th>Parameter</th>
<th>Value</th>
</tr>
</thead>
<tbody>
<tr>
<td>Bunch population $N_p$</td>
<td>$1 \times 10^{11}$</td>
</tr>
<tr>
<td>Bunches</td>
<td>20</td>
</tr>
<tr>
<td>Bunch spacing $L_{sep}$</td>
<td>7.48 m</td>
</tr>
<tr>
<td>Bunch profile</td>
<td>Gaussian</td>
</tr>
<tr>
<td>Rms bunch length $\sigma_z$</td>
<td>7.7 cm</td>
</tr>
<tr>
<td>Rms transverse size $\sigma_{x(y)}$</td>
<td>300 $\mu$m</td>
</tr>
<tr>
<td>Chamber radius $r$</td>
<td>2 cm</td>
</tr>
<tr>
<td>Magnetic field $B$</td>
<td>0 T</td>
</tr>
<tr>
<td>Primary $e^-$ rate</td>
<td>$10^{-3}$ e$^-$/p$m$</td>
</tr>
<tr>
<td>Max SEY $\delta_{\text{max}}$</td>
<td>2.03</td>
</tr>
<tr>
<td>Energy at max SEY $E_{\text{max}}$</td>
<td>262 eV</td>
</tr>
</tbody>
</table>

**Figure 2: POSINST SEY model for Copper, with $\delta_{\text{max}}=2.03$, $E_{\text{max}}=262$ eV, for normal incidence.**

Figure 3: Buildup results for a) the ECLOUD SEY model (in blue) or b) the POSINST-like SEY model in full (in red) or with some components suppressed (magenta and green).
Figure 4: Energy distribution curves of secondary electrons as a function of the incident energy, as obtained with the ECLOUD (blue curve) or the POSINST-like (red curve) model.

Figure 5: Experimental data from surface measurements recently performed at CERN [3]. Plotted are SEY curves for different scrubbing conditions.

is shown in Fig. 4 and is especially evident at low incident energies (the total spectrum of the true secondaries is in both cases shaped according to Eq. 6).

**EXTRACTING A THIRD MODEL FROM RECENT SEY MEASUREMENTS: FULL SCRUBBING CASE**

Recent experimental measurements of secondary emission from a Copper surface were obtained at CERN on a prototype of the LHC beam screen under different scrubbing conditions [3]. Secondary yields were measured (see Fig. 5), as well as related energy distribution curves (EDC) of the secondary electrons as a function of the primary electron energy, and special attention was paid to measure-

ments for very low-energy electrons (≤20 eV).

A first study was made on the data for the full scrubbing case (δmax ≃1), trying to fit empirical curves to the yield measurements, and testing different models for the secondaries’ energy spectrum to find which one reproduced the observations most closely. Fig. 6 shows the curves used to describe the yield as a function of incident energy: a function similar to Eq. 2 was used to model the true secondary part (green line), whereas the elastic component (in blue) was obtained by subtraction from the data for the total yield (in red). Two different distributions were studied to model the total secondary electrons’ energy spectrum: the one proposed in [1] and described by Eq. 6, and an alternative Lorentzian model, originally proposed in [11] and slightly modified to constrain it to drop to zero at zero incident energy. The resulting electron spectra are shown in Fig. 7. Results of the study are shown in Fig. 8, where EDC curves obtained at different primary incident energies

Figure 6: Fully scrubbed Copper case: SEY measurements are shown in red, and the fitted true secondary and elastic components in blue and green respectively.

Figure 7: Secondary electrons’ energy spectra, when using either a) the distribution defined in Eq. 6 (blue line), or b) a quasi-Lorentzian model (red line).
Figure 8: EDC curves for different incident energies: in black are the measurements from [3], in blue the simulation results for the case in which model a) of Fig. 7 was adopted, and in red the results for the quasi-Lorentzian model b).

Figure 9: Comparison of build-up results with the two SEY models described for the reference case of Table 1.

for the two models adopted are compared with measurements published in [3]. The difference between the two models is not very pronounced, although for higher primary energies the Lorentzian approach gives a better description of the lower energy part of the spectrum. Also some improvements could be introduced in both SEY models to better reproduce the data in the intermediate energy spectrum. Simulation results for the electron cloud build-up in a representative field-free section of the chamber for the reference case of Table 1, when using the yield curves for the full scrubbing case and the two models for the secondaries’ energy spectrum, are shown in Fig. 9.

Figure 10: SEY measurements for an ‘as received’ Copper surface (in red) and a few fit attempts (in blue, green and magenta). In the top right inset a zoomed-in view of the high energy tails (where no data is available) is given.

FITTING DATA FOR ‘AS RECEIVED’ SURFACE CASE (NO SCRUBBING)

The same study has been repeated for the set of SEY measurements obtained for an ‘as received’ Copper surface. This case, however, is complicated by the fact that, at higher secondary emission yields (up to $\delta_{\text{max}} \simeq 2$) that enhance the electron multipacting, even small differences in the curves fitting the data can lead to substantial effects on the build-up results. In particular, a high level of uncertainty is present in the definition of the curves at high incident energy, since data are provided only in the range 0-350 eV, but in the reference case examined electrons can be accelerated up to 1.5 keV. Three different fits have been attempted here, as shown in Fig. 10: the true secondary component has been modelled via a function like Eq. 2 weighted by a factor $e^{-\pi/p}$ (where $p$ is a fitting parameter), and the elastic component has been again obtained by subtraction.

Build-up results for these three different fits when using either a quasi-Lorentzian model for the secondaries’ energy spectrum or the distribution from Eq. 6 are given in Fig. 11 and 12 respectively. There is a difference of up to a factor of 3 in the level of electron cloud line density reached with the three different fits, and the choice of distribution to model the secondary energy spectrum alone accounts for a difference in the results by approximately a factor of 2. A broader picture can be obtained from the scan in Fig. 13, where the accumulated peak electron line density is plotted as a function of the bunch intensity for the six cases studied above. Even if the results are strongly sensitive to the details of the secondary emission model used, the build-up pattern is fairly constant and the intensity threshold is not severely affected. Comparing with the results presented in Fig. 3 for approximately the same $\delta_{\text{max}} \simeq 2$, it is possible to see how the values of the electron cloud peak line density...
CONCLUSIONS

We have presented results of a recent study exploring the sensitivity of the electron cloud buildup mechanism to the details of the secondary electron emission model adopted in the simulations. A tentative phenomenological fit of recent SEY experimental measurements on Copper has been described, the resulting parametrisations tested in simulations and the results compared to existing models. Electron cloud build-up results show a strong sensitivity to the models’ details (qualitative and quantitative separation of the SEY components and total energy spectrum assigned to the secondary electrons), though the intensity threshold level is not severely affected.

A level of uncertainty is however present in the fit precision at high energy, due to the lack of experimental data in the region. In conclusion, more measurements of SEY properties at primary incident energies $E_0 \geq 400$ eV and EDC data are needed to achieve a better precision in the extraction of a phenomenological model.

ACKNOWLEDGEMENTS

The author would like to gratefully acknowledge R. Cimino for kindly providing all experimental data for this study and for related discussions; G. Rumolo for his collaboration and helpful support; F. Zimmermann, M. Furman and M. Pivi for valuable communications.

REFERENCES


Electron Cloud Instability Simulations for the LHC

E. Benedetto*, F. Zimmermann, CERN AB/ABP, Geneva, Switzerland

Abstract

We present results of simulations of transverse single bunch instabilities for the LHC and LHC Upgrade scenarios and in particular first results with electron cloud in the dipole regions. The possibility to use the code HEADTAIL below the threshold of the fast instability is under study and we here discuss a possible way to distinguish numerical noise from true physics.

INTRODUCTION

The instabilities induced by electron cloud are a concern for LHC and for the future upgrades of the collider, especially at injection energy. Simulations have been done with HEADTAIL [1, 2, 3]. The code models transverse, single bunch instabilities, assuming that the interaction between the electron cloud and the bunch happens at a finite number of locations around the ring. The electrons and the protons are represented by macroparticles and the bunch is divided into slices. A PIC module compute the interaction between the 2D cloud and each slice. In the following section we show first results of modelling the real distribution of the electrons in dipole field regions. Then some prediction for LHC upgrade scenarios are discussed. Finally, the last section is devoted to the study of the emittance growth below the threshold of the fast head-tail instability and its dependance on the number of interaction points per turn.

ELECTRON CLOUD IN THE DIPOLES

It has been observed that in the SPS the electron cloud is mainly concentrated in the dipole regions. The electrons populate stripes at a certain distance from the beam, depending on the bunch intensity. We suppose that this will be valid also for the LHC and, since the nominal bunch intensity will be the same, we can assume (from SPS measurement [4]) that the stripes will be at about 14 mm from the axis, with an rms size of $\sigma_{str} = 8$ mm. In these simulations, 10% of the electrons are distributed uniformly inside the chamber and the other 90% are populating the stripes, which have a uniform profile, $4\sigma_{str}$ wide. Figure 1 shows the projection of the electron cloud density onto the horizontal axis (the chamber is elliptical) and the electron cloud density evolution in the vertical plane. It is evident that the electrons pinch toward the center during the passage of the bunch in the vertical plane, while in the horizontal plane they do not move because of the strong magnetic field approximation, which freezes their motion perpendicular to the field lines.

The emittance growth (Fig. 2) is very low and even for high average cloud density we are below the threshold of the fast head-tail instability. If the electrons are really concentrated in stripes, far from the beam axis, an average density higher than the one predicted for uniform distribution seems to acceptable, even at injection. However there are a lot uncertainties in the electron distribution and, moreover, the stripes position is not fixed, but varies with the beam parameters.

ELECTRON CLOUD SIMULATIONS FOR THE LHC UPGRADE

Simulations for LHC upgrade scenarios have been done pessimistically assuming that the electron cloud is uniformly distributed in the transverse plane (no magnetic field and no stripes).

In Fig. 3 the emittance growth vs. time is shown for the different scenarios, whose parameters are listed in Table 1. The electron cloud density is $\rho_e = 1.2 \times 10^{12}$ m$^{-3}$. Even...
Table 1: Parameters for the nominal LHC and for LHC Upgrade Scenarios

<table>
<thead>
<tr>
<th>parameter</th>
<th>nominal LHC (ultimate)</th>
<th>Short bunch</th>
<th>Piwinski</th>
</tr>
</thead>
<tbody>
<tr>
<td>cloud density, (\rho_e [m^{-3}])</td>
<td>(1.2 \times 10^{12})</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>bunch population, (N_b)</td>
<td>(1.1 \times 10^{11}(1.7 \times 10^{11}))</td>
<td>(1.7 \times 10^{11})</td>
<td>(3.0 \times 10^{11})</td>
</tr>
<tr>
<td>beta function, (\beta_{x,y} [m])</td>
<td>100</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>rms beam size, (\sigma_{x,y} [mm])</td>
<td>0.884</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>rms bunch length, (\sigma_z [m])</td>
<td>0.13</td>
<td>0.0806</td>
<td>0.214 (uniform profile)</td>
</tr>
<tr>
<td>rms momentum spread, (\delta)</td>
<td>(4.68 \times 10^{-4})</td>
<td>(7.71 \times 10^{-4})</td>
<td>(1.40 \times 10^{-4})</td>
</tr>
<tr>
<td>longitudinal emittance, (\epsilon_z [eVs])</td>
<td>1.25</td>
<td>1.25</td>
<td>0.625</td>
</tr>
<tr>
<td>synchrotron tune, (Q_s)</td>
<td>0.0059</td>
<td>0.0188</td>
<td>0.00128</td>
</tr>
<tr>
<td>momentum compact fact, (\alpha_c)</td>
<td>(3.47 \times 10^{-4})</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>circumference, (C [m])</td>
<td>26659</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>nominal tunes, (Q_{x,y})</td>
<td>64.28, 59.31</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>chromaticity, (Q'_{x,y})</td>
<td>2.2</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>dispersion, (D [m])</td>
<td>1.5</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>magnetic field</td>
<td>no</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>relativistic factor, (\gamma)</td>
<td>479.6</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>cavity voltage, (V [MV])</td>
<td>8</td>
<td>21.5</td>
<td>3</td>
</tr>
<tr>
<td>harmonic number, (h)</td>
<td>35640</td>
<td>106920</td>
<td>3564</td>
</tr>
</tbody>
</table>

Figure 2: Emittance growth vs. time, with different average electron cloud densities, for LHC at injection for the nominal LHC a strong instability would develop, if the cloud distribution was uniform. Assuming this electron cloud density for all cases, we can compare the different scenarios and it seems that the “short bunch” option would be the most favorable, due to the fact that the synchrotron motion mixes head and tail protons more quickly. On the other side, the “Superbunch”, or “Piwinski” configuration, is the most unstable one (here ignoring that the real electron density may be much lower for this case).

SLOW TERM EMITTANCE GROWTH

Figure 2 shows that even below the threshold of the fast head-tail instability, a slow, long-term emittance growth is still present, which depends on the electron density and seems to be linear in the time. Studies are ongoing to understand how much the simulation parameters (number of macroparticles, grid size, number of points of interaction between the beam and the cloud) influence the results in this regime [5, 6]. We here discuss the dependance on the number of kicks per turn. A scan on the number of interaction points has been done, assuming the nominal LHC parameters, and a low electron density of \(\rho_e = 2 \times 10^{11} m^{-3}\), in order to stay below the fast headtail instability threshold. Figure 4 shows the relative emittance growth rate vs. the number of interaction points along the ring. No clear convergence is seen. This is probably due to the fact that when we set the number of kicks per turn, we change the effective phase advance between the interaction points and the strength of the kick itself. The plot (Fig. 5) of the emittance growth rate (multiplied by the number of kicks) vs. the fractional part of the tune divided by the number of kicks shows the possibility of hitting some resonances.

Figure 3: Vertical emittance growth vs. time for the different LHC scenarios, at injection energy.
A tune scan was done, with one kick per turn, for an electron cloud density of \( \rho_e = 4 \times 10^{10} \text{ m}^{-3} \) in order to stay in the slow emittance growth regime. Figure 6 clearly reveals the presence of a resonance pattern. Anyway, understanding which resonance is responsible for a larger emittance growth rate is quite difficult. In Fig. 7 the scanned working points are superimposed on the resonance tune diagram. Shown in red are the points which correspond to larger emittance and in black some of the low-emittance ones. It is not clearly evident whether these points lie on resonance lines or not. In addiction, due to the interaction with the electron cloud, the working point can change. In orange, we plot the ‘effective tune’ corresponding to cases of larger emittance growth, computed via an FFT of the bunch centroid. Only in one case, corresponding to the third order resonance line \( Q_x = 0.666 \), the nominal and the effective point agree, but normally the effective tune is shifted. Figure 8 shows an example of the FFT of the horizontal and vertical bunch centroid motion, from which the effective tune was obtained.

**CONCLUSIONS AND OPEN QUESTIONS**

The electrons will presumably be concentrated in the dipole field regions of the LHC ring and they will mainly populate stripes, whose distance from the axis depends on the bunch intensity. If they are far from the beam, the induced emittance growth will be lower than for the uniform distribution. First results for the LHC underline the need to take into account and model the real electron distribution around the ring. Benchmarking with experiments in the SPS is in progress. The simulations for LHC upgrade scenarios have been done assuming a constant electron cloud distribution. This is quite a pessimistic assumption, but it gives an idea of differences between the scenarios. Studies are ongoing on the possibility to use HEADTAIL code.
below the threshold of the fast instability and to distinguish numeric noise from true physics. We should be aware that some effects in the emittance growth are due to hitting some resonance excited by the placement of the interaction points. The problem raises whether this is just an artifact of the simulations or whether we have to take into account the cloud modulation along the ring. Once more, we feel the need to model the real distribution of the electrons and for this purpose in a collaboration with USC, we are implementing in the quasi-continuous plasma code QuickPIC the characteristics of both the beam and the electron distribution for the individual magnets elements of the ring. In this way we suppress artificial resonances and may observe true ones.

ACKNOWLEDGEMENTS

The authors thank G. Arduini, A. Ghalam, T. Katsouleas, E. Metral, K. Ohmi, F. Ruggiero, G. Rumolo and D. Schulte for helpful comments and discussions.

REFERENCES

COHERENT RADIATION OF ELECTRON CLOUD

S. Heifets
Stanford Linear Accelerator Center, Stanford University, Stanford, CA 94309, USA

1

Abstract

The electron cloud in positron storage rings is pinched when a bunch passes by. For short bunches, the radiation due to acceleration of electrons of the cloud is coherent. Detection of such radiation can be used to measure the density of the cloud. The estimate of the power and the time structure of the radiated signal is given in this paper.

2 INTRODUCTION

Effect of the electron cloud on the transverse beam stability in the storage rings has been discovered in the KEK Photon Factory[1] and studied extensively afterwards [2], [3]. The instability is one of the culprits limiting performance of the B-factories. The density of the cloud is the main parameter defining the instability. Diagnostics of the density could be important in limiting adverse effects of the cloud. Here we study the coherent radiation from the cloud as a possible tool for diagnostics of the e-cloud density. The radiation is the result of acceleration of electrons in transverse to the beam plane by the field of passing bunches. Radiation is coherent and propagates in the beam pipe. It can be detected with an rf antenna. The amplitude of the signal depends on the density of the cloud and can be used for the cloud diagnostics. We estimate here the power and the time structure of the signal.

3 MODEL

The beam is a train of the high energy positron bunches. Electron cloud is generated by synchrotron radiation and/or multipaction of secondary electrons. Electrons are accelerated radially by the field of a passing bunch and radiate during the passage of a bunch. Radiation field can be expanded over the eigen-modes $E_{m,n}^{(\pm)}$, $H_{m,n}^{(\pm)}$ in the beam pipe. Consider a round beam pipe with the radius $b$. The typical $b$ are of few cm and the cut-off frequency of the beam pipe is several GHz. Let us assume the frequency dependence in the form $e^{-\omega t}$, use the cylindrical coordinate system with the axis along the beam line, and denote by $z_d$ the point of observation of the radiation. The Fourier harmonics of the magnetic field radiated at the frequency $\omega$ by the current induced in the electron cloud and propagating in $\pm z$ directions are

$$H^{\pm} = \sum_{m,n} a_n^{(\pm)} H_{m,n}^{(\pm)} e^{i\phi_{m,n} (z-z_d)},$$

and similar equation for $E_{m,n}^{(\pm)}$. Here $m$ is the azimuthal number, and $n$ the radial mode number. The propagating constant $q_{m,n} = \frac{k^2 - \lambda_{m,n}}{\lambda_{m,n}}$ is defined with the cut in the complex plane of $\omega$ along the real axis from $k = -\lambda_{m,n}$ to $k = \lambda_{m,n}$. On the upper edge of the cut $q_{m,n} = +i\sqrt{\lambda_{m,n}^2 - k^2}$. The small positive imaginary part is implied, $k = \omega/c + i\epsilon$, $\epsilon = 0+$, and $q_{m,n}(-k^*) = -q_{m,n}^*(k)$.

For TM mode, $\lambda_{m,n} = \nu_m / b$, where $\nu_m$ is the $n$-th root of the Bessel function $J_m(\nu_m) = 0$.

The explicit form of the eigen-modes can be found in the text books. In the following, we assume that the induced current in the cloud is axially symmetric and has only radial component. In this case, it is suffice to consider only $m = 0$ mode, and we drop this index below. The nonzero components of azimuthal $m = 0$ TM modes are

$$E^{s,r}_n = i\frac{q_n}{\lambda_n} J_n(\lambda_n r) e^{i q_n (z-z_d)},$$
$$E^{s,z}_n = J_0(\lambda_n r) e^{i q_n (z-z_d)},$$
$$H^{s,\phi}_n = \frac{i}{\lambda_n} J'_0(\lambda_n r) e^{i q_n (z-z_d)},$$

where prime means derivative over the argument, $s = \pm 1$, $\lambda_n = \nu_n / b$, and $\nu_n$ is the $n$-th root of $J_0(\nu_n) = 0$.

The modes are orthogonal with the norm $N_{n,m}^{s,\phi}$ being proportional to the integral of the Pointing vector over the cross section of the beam pipe,

$$N_{n,m}^{s,\phi} = \int S(E^{s,r}_n H^{s,\phi}_m - E^{s,\phi}_n H^{s,r}_m) = 2\pi b^2 \frac{k q_n}{\lambda_n} J_n^2(\nu_n) \delta_{n,m} \delta_{s,-s'}.$$ (3)

The amplitudes of the modes can be found from the identity following from Maxwell equations for the free field $E_2$, $H_2$ and the fields $E$, $H$ driven in the same volume $V$ by the current $j$,

$$\int dS [E_1 \times H_2 - E_2 \times H_1] = -Z_0 \int dV E_1 j_2.$$ (4)

Here $Z_0 = 4\pi/\epsilon_0 = 120\pi$ Ohms.

Substituting $E$, and $H$ from Eq. 1 we get

$$a_n^{\pm} = \pm \frac{Z_0}{\lambda_{m,n} \delta_{n,m}} \int dV' j'_n(r',z') E_n^{\mp,r}(r',z').$$ (5)

The integral in $a_n^{\pm}$ over $z$ is taken over the region where the modes were generated: $-\infty < z < z_d$ for $a_n^+$ and $z_d < z < \infty$ for $a_n^-$. 271
The relativistic bunch gives the radial kick to an electron of the cloud. Therefore, in the axially symmetric structures, the current has only radial component.

The current can be defined from the equation of motion for the trajectory \( R(r, z, t) \)

\[
\frac{d^2 R}{dt^2} = \frac{2N_0 r_e^2 R}{R^2 + \sigma_\perp^2} \rho_B(z - ct),
\]

where \( r_e \) is classical electron radius, \( N_0 \) is bunch population, and \( \rho_B(z) \) is the longitudinal density profile, \( \int \rho_B(z)dz = 1 \). We set the initial conditions at \( t = -\infty \)
\n\[
R(r_0, z, t \to -\infty) = r_0, \frac{\partial R(r_0, z, t \to -\infty)}{\partial t} = 0.
\]

It follows from Eq. 6 that \( R(r_0, z, t) \) depends only on the initial location \( r_0 \) and the parameter \( \zeta = ct - z \). Let us assume the uniform density of the cloud, \( n(r_0, z_0) = n_0 \). Then, the current gives

\[
j(r, z, t) = \int 2\pi n_0 dr_0 dz_0 n(r_0, z_0) \frac{\partial}{\partial t} R(r_0, z_0, t)
\]

\[
\delta(z - z_0) \frac{\partial}{\partial t} R(r_0, z_0, t)
\]

\[
e^{\frac{n_0 r_e}{r}} \int \frac{r_0 dr_0 R(r_0, \zeta)}{\lambda_n R(r_0, \zeta)} \delta[r - R(r_0, \zeta)]
\]

is also a function of \( r \) and \( \zeta = ct - z \), \( j(r, z, t) = j(r, \zeta) \). The integral in Eq. 5 over \( dr' \) gives

\[
\int r' dr' j_0^\lambda(\lambda_n r') j(r', z', t') = e^{\frac{n_0 r_e}{r}} \int \frac{dr'_0 d\lambda'_0}{\lambda'_0 \lambda_n} \frac{dR(r_0, \zeta')}{d\zeta'}
\]

Changing here integration over \( t' \) to integration over \( d\zeta' \), \( \zeta' = ct' - z' \), and substituting result in Eq. 5 we can carry out integration over \( dz' \). That gives

\[
e_{n}^{(z)} = \frac{4\pi n_0 r_e^2}{\lambda_n} \int d\zeta e^{i k z_a} \frac{\partial f(\zeta)}{\partial \zeta}
\]

where

\[
f(\zeta) = \int_0^b \frac{r_0 dr_0 J_0(\lambda_n R(r_0, \zeta))}{\lambda_n R(r_0, \zeta)}
\]

Let us assume that the amplitude of a signal in the detector situated at \( z = z_d \) is determined by the component \( E_d(r = b, z_d, t) \). The field at \( z = z_d \) is

\[
E_{r}(r, z_d, t) = \frac{1}{2} \sum_{m} [E_{m}^{(+)r} a_{m}^{(+)} r_{d}(r, z_d) + a_{m}^{(-)} E_{m}^{(-)r}(r, z_d)],
\]

and the form

\[
E_{r}(r, z_d, t) = \theta(ct - z_d) \sum_{m} \frac{4\pi n_0 J_1(\nu r_e)}{\lambda_n^2} \frac{e^{i k z_a}}{b^2 \Gamma_1} \frac{d^2 f(\zeta)}{d\zeta^2} + \lambda_n^2 f(\zeta_d) - \lambda_n^2 f(-\infty)).
\]

Here \( \zeta_d = ct - z_d \), and the last term \( f(-\infty) = \) \( \sum_{m} \frac{4\pi n_0 J_1(\nu r_e)}{\lambda_n^2} \frac{e^{i k z_a}}{b^2 \Gamma_1} \frac{d^2 f(\zeta)}{d\zeta^2} + \lambda_n^2 f(\zeta_d) - \lambda_n^2 f(-\infty)) \).

In deriving Eq. 12 we used the integral

\[
\int_{-\infty}^{\infty} \frac{dk}{k} \frac{q_n(k)}{k + q_n} - \frac{q_n(k)}{k - q_n} \frac{e^{i k z_a}}{k} = \frac{2}{\lambda_n^2} \int_{-\infty}^{\infty} \frac{dk}{k} \frac{\lambda_n^2(k)}{k} \frac{e^{i k z_a}}{k},
\]

\[
= 2\pi i \left[ -\frac{\partial}{\partial \zeta} \delta(z_d - ct + \zeta') + \lambda_n^2 \theta(z_d - ct + \zeta') \right],
\]

where \( \theta(x) \) is the step function.

The power radiated by bunch per turn is \( P = \Delta U/T_0 \) where \( T_0 \) is the revolution period and \( \Delta U \) is radiated energy.

\[
\Delta U = \frac{e}{4\pi} \int \frac{d\omega}{2\pi} |\mu_n|^2 |N_n n_n|.
\]

\( \Delta U \) is given by the sum over modes propagating in the beam pipe. The sum converge very rapidly and the main contribution is given by the lowest mode.

To proceed further, we consider two extreme cases of very short and very long parabolic bunches.

In the first case, interaction with the bunch gives instantaneous kick. The trajectory given by Eq. 6

\[
R(r_0, \zeta) = r_0 - \frac{2N_0 r_e r_0}{\sigma_\perp + r_0} \zeta,
\]

defines the function Eq. 10, \( f(\zeta) = \nu^2 F(\zeta/\Delta) \), where

\[
\frac{\partial}{\partial \zeta} \delta(z_d - ct + \zeta') + \lambda_n^2 \theta(z_d - ct + \zeta') \right],
\]

and

\[
1 + \frac{2N_0 r_e}{b^2} = \frac{e Z_0 I_{bunch} \mu_e}{mc^2} \frac{R}{b^2},
\]

where \( mc^2 \) is electron mass.

The field at \( r = b \) takes the form

\[
E_{r}(r, z_d, t) = \Delta \theta(\zeta) \Phi(\zeta/\Delta),
\]

and \( \Phi \) is given by Eq. 26.
\[ \Phi(\xi) = \frac{d^2F(\xi)}{d\xi^2} + \left( \frac{\nu L}{b} \right)^2 F(\xi) - \nu \left( \frac{L}{b} \right)^2 J_1(\nu), \]  

(20)

where \( F \) is related to \( f, f = \nu^2 F \).

Because the sum converge very rapidly, we retain only the first term in the sum \( \nu = \nu_1 \approx 2.4 \).

With the PEP-II parameters \( \beta_{\text{bunch}} = 2.5 \text{ mA}, 2\pi R = 2 \text{ km}, b = 2.5 \text{ cm}, \) and assuming typical \( n_0 = 10^9 \text{ cm}^{-3}, \) we get \( \Delta = 1.91 \times 10^3 (e/b)^2 \). The function \( \Phi(\xi) \) is shown in Fig. 1. It gives the shape of the signal measured at the wall for a short bunch provided the signal is proportional to \( E_r \). For short bunches, the signal is due to radiation of the instantaneously accelerated electrons and variation in time is defined by the parameter \( \gamma_0 \).

Figure 1: Time dependence of a signal proportional to \( E_r \) at the wall for short bunches.

For parabolic bunches with the length \( 2L \)

\[ \rho(\zeta) = \frac{3}{4L} \left( 1 - \frac{\zeta^2}{L^2} \right), \]

(21)

and electrons within the beam, \( r_0 < \sigma \), the equation of motion describes oscillations of the trapped electrons,

\[ \frac{d^2R}{d\xi^2} + \Omega^2(\xi) R = 0, \]

\[ \Omega^2(\xi) = \frac{2Nv_\rho L^2}{\sigma^2} \rho(\xi) = \Omega_0^2(1 - \xi^2), \]

(22)

where \( \xi = \zeta/L, \zeta = ct - z \) and

\[ \Omega_0^2 = \frac{3Nv_\rho L}{2\sigma^2}. \]

(23)

The field in this case is given by

\[ E_r(b, z, t) = \Lambda \theta(\xi) \Phi(\xi), \]

(24)

where

\[ \Lambda = \frac{4\pi en_0 b^3}{\nu^3 L^2 J_1(\nu)}, \]

(25)

\[ \xi = \zeta/L \] and \( \Phi \) is given by Eq. 26.

\[ \Phi(\xi) = \frac{d^2F(\xi)}{d\xi^2} + \left( \frac{\nu \gamma_0}{b} \right)^2 F(\xi) - \nu \left( \frac{\gamma_0}{b} \right)^2 J_1(\nu). \]

(26)

With the same parameters for \( n_0 \) and \( b, \) and for \( L = 10 \text{ cm}, N_0 = 10^{13}, \sigma = 5 \text{ mm}, \) we get \( \Delta = (e/b)^2 \times 6.8 \times 10^5. \) The time dependence in this case is shown in Fig. 2 and is defined by the frequency \( \Omega_0 \).

Figure 2: Time dependence of a signal proportional to \( E_r \) at the wall for long parabolic bunches.

Let us estimate the energy radiated by a long bunch. The radiated energy Eq. 14 is

\[ \Delta U = \frac{4\pi}{\nu^2} \left( \frac{e\gamma_0}{J_1(\nu)} \right)^2 \int dk \frac{1}{k} \frac{1}{k - q} | \int d\zeta e^{ik\zeta} \frac{df}{d\zeta} |^2. \]

(27)

Here we neglected contribution of all terms except \( n = 1 \), and denote \( \nu = \nu_1, q = \sqrt{k^2 - (\nu/b)^2}. \)

The radiation is due to oscillations of electrons trapped in the bunch. Therefore, we can assume that for such electrons \( R(r_0, \zeta) \approx \sigma \), \( \zeta < b \). Because significant contribution comes from the first Bessel root \( \nu_1 \approx 1, \) that allows us expand \( J_0(\lambda_\nu r_0) \approx 1 - \lambda_\nu r_0^2/2 \). The trajectory with the initial conditions \( R(r_0, 0) = r_0, \) and \( (dR/d\zeta)|_{\zeta=0} = 0, \) is \( R(r_0, \zeta) = r_0 \cos \psi(\zeta) \). In this approximation,

\[ f(\zeta) = \frac{\nu^2}{2} - \left( \frac{\lambda_\nu b^2}{4} \right)^2 \cos^2 \psi(\zeta), \quad \zeta > 0; \]

\[ \int d\zeta e^{ik\zeta} \frac{df}{d\zeta} = \left( \frac{\lambda_\nu b^2}{4} \right)^2 \int d\zeta \Omega(\zeta) e^{ik\zeta} \sin(2\psi(\zeta)). \]

The last integral can be evaluated by the saddle-point method. The integral is exponentially small for \( 2\Omega_0/(kL) < 1 \), and

\[ \int d\zeta \Omega(\zeta) e^{ik\zeta} \sin(2\psi(\zeta))^2 \approx \left( \frac{\nu b}{4} \right)^2 \frac{\Omega_0^2}{\sqrt{(2\lambda_\nu/kL)^2 - 1}}. \]

(29)
otherwise. The radiated energy is given by $\Delta U = \frac{e^2}{b} \frac{1}{2\nu L} \left( \frac{\pi \nu_0 b^2}{4} \right) \left( \frac{\nu}{\nu_c} \right)^2 \left( \frac{L}{b} \right) S \left( \frac{2\Omega_0 b}{\nu L} \right)$, \( \times 10^{13} \), the typical density of the cloud

$$S(p) = \int_1^p \frac{x^2 dx}{(x-\sqrt{x^2-1})^2 \sqrt{p^2-x^2}}$$

where

The function $S(p)$ grows fast with $p$ as $p^4$, see Fig. 3.

Taking $N_{\text{bunch}} = 10^{13}$, the typical density of the cloud $n_0 = 10^3 \text{ cm}^{-3}$, $\sigma_\perp = 0.5 \text{ cm}$, $L = 10 \text{ cm}$, and $b = 2.5 \text{ cm}$ we get $\Omega_0 = 12.9$, $p = 2.7$, $S(p) = 181.7$, and $\Delta U = 3.17 \times 10^{10} \left( \frac{e^2}{b} \right)$, about $0.7 \mu J$.

Figure 3: Function $S(p)$.

4 SUMMARY AND DISCUSSION

We give the estimate of the power radiated by the pinched electron cloud after passage of a positron bunch. The power seems to be detectable and the detection of the signal can provide information on the density of the cloud. The main difficulty, probably, is separation of the signal from the signal induced by the bunch and from the wake field induced by geometric discontinuities and the noise of the cloud. Placement of the detector in a straight pipe may help to suppress the wake fields. Dependence of the signal on current $P \propto n_0^2$ is quite different from the linear dependence of a signal on current due to wakefields because the electron cloud density in saturation $n_0$ itself proportional to beam current. The frequency content of the signal depends on the bunch length and on the transverse distribution of electrons in the cloud: electrons in the vicinity of the beam oscillate with plasma frequency $(\omega/e)^2 = 4\pi m_e$ while electrons with initial radial position $r(0) \gg \sigma_r$ shift slightly during the kick from the bunch. That also distinguish the signal from the signal due to the long-range wake fields which have a narrow bandwidth around the HOM frequency. The time structure of the signal is different from the bunch profile. It would be interesting to study effect on the signal of the solenoidal magnetic field. Such a field generated in the beam pipe to suppress the cloud changes the dynamics of the electrons and would affect the signal.

It is worth mentioning the attempt to measure the density of the cloud by detecting the phase shift of the RF wave induced in the beam pipe [4]. Unexpectedly, it was discovered that the passing bunch strongly affects the detected signal. We suggest that the interference may be explained by the coherent radiation of the cloud described above.

Acknowledgments

I thank M. Zolotorev and the theory ARDA group for helpful discussion. Work supported by Department of Energy contract DE–AC03–76SF00515.

5 REFERENCES

Abstract
This paper provides a summary of Session 1 of the workshop. The session laid the basis for the subsequent detailed presentations and discussions on options for increasing the luminosity of the LHC. In particular, the session summarised the physics motivation and experiment considerations, sketched the critical items required for improvements to the LHC luminosity, outlined options for upgrading the LHC insertion regions and injectors, and introduced the technological constraints and challenges for such an upgrade.

INTRODUCTION TO CERN FUTURE ACTIVITIES
In his opening talk, the CERN Director-General, R. Aymar, outlined the laboratory’s mission and possible directions for future research. He stressed that while the highest priority will remain the completion of the LHC accelerator and experiments and the start of LHC operations in summer 2007, the proposed strategic orientations for CERN in the years 2004-2010 include the development of detailed technical solutions for a future upgrade to the LHC luminosity, to be commissioned in the years 2012-2015.

In line with the new policy of the European Commission for re-structuring the European research area, it was proposed to launch in the years 2004-2006 a range of studies for upgrading the LHC luminosity in cooperation with other laboratories. The studies should be along the following lines:

- Design of the 160 MeV proton LINAC through the European Programme for a High Intensity Pulsed Proton Injector (HIPPI).
- Modifications to the magnets in the interaction regions at two crossing points of the LHC beams, linked with the European programme Next European Dipole (NED), which aims to construct superconducting magnets operating up to 15 T.
- Definition of the required upgrades to the ATLAS and CMS detectors in order to withstand a factor of 10 higher luminosity.

A review shall take place in 2009-2010 to re-define the overall strategy for CERN activities for the decade 2011-2020, in light of the first physics results from the LHC and of progress with previous R&D efforts and actions. Possible choices are presently quite open.

The CERN Chief Scientific Officer, J. Engelen, underlined the importance of new multi-GeV high intensity proton accelerators. The Superconducting Proton Linac (SPL) is proposed to deliver 4 MW of beam power at up to 2.2 GeV energy while rapid cycling synchrotrons delivering the same beam power but up to 50 GeV energy have been proposed as a complement to the SPL option. Such new accelerators will allow higher intensities to be achieved throughout the CERN complex, thus being of great utility for both the high luminosity operation of the LHC and for the CERN fixed target physics experiments.

PHYSICS MOTIVATION
An upgrade to the LHC luminosity will provide an extension to the LHC physics discovery reach and improve precision measurements of parameters. Examples of such gains are given below.

Figure 1 shows the improved reach for Higgs bosons within the Minimal Supersymmetric Standard Model (MSSM). Should a (light) Higgs not be observed at the LHC, an upgraded LHC could be used to study a new strong interaction regime in $V_fV_f$ scattering.

A higher integrated luminosity also brings an increase in the supersymmetric (SUSY) particle discovery potential. Increasing the integrated luminosity from 300 fb$^{-1}$ (LHC) to 3000 fb$^{-1}$ (LHC upgrade) will increase the squark and gluino mass reach by ~500 GeV to 3 TeV.
The increased statistics of an upgraded LHC would also be used to improve precision measurements. Studies of Higgs couplings, triple gauge boson couplings and mass measurements of SUSY particles would all stand to benefit from an increased luminosity. For all these physics studies, the total integrated luminosity, acquired in stable machine running conditions, should not be compromised by an increase in the instantaneous luminosity.

**DETECTOR CONSIDERATIONS**

To be able to profit fully from an increased luminosity, the detector performance at an upgraded LHC should be as for the LHC. To maintain the same performance as the LHC, the Trackers need to be replaced by new detectors based on improved radiation-hard sensors and electronics. Although for the most part the calorimeters and muon spectrometers do not need to be changed, the energy resolution of the former and the occupancy of the latter will be degraded due to the additional hits coming from pile-up events. Scintillators and certain end-cap electronics will, however, need to be exchanged. The Level-1 trigger electronics and processors will need to be replaced in the case of a change in the accelerator bunch crossing frequency.

It was noted that discoveries at high masses could be performed using information only from the calorimeters and the muon systems while precision measurements and the understanding of new phenomena require the identification and precise reconstruction of high $p_T$ objects (electrons, photons, $\tau$-leptons, b-jet), good tracking capabilities and background rejection and improved forward jet tagging.

The number of pile-up inelastic events per bunch crossing is about 20 for the nominal LHC scheme of 25 ns bunch spacing. In the case of the Super-bunch option (see Table 1 below) and assuming no changes to the trigger scheme, the corresponding number rises to 25000. It would be extremely difficult for the experiments to exploit the increased luminosity under these conditions. In order to reduce the effect of pile-up events, and thus easing pattern recognition in the trackers and calorimeters, the experiments prefer shorter but finite bunch spacing (e.g. 10, 12.5 or 15 ns.)

**POSSIBLE SCENARIOS FOR AN LHC LUMINOSITY UPGRADE**

As Figure 2 shows, due to the high radiation doses, the life expectancy of the LHC insertion region quadrupoles is estimated to be less than 10 years. Therefore, it is reasonable to plan for a machine luminosity upgrade based on new low-$\beta$ insertion region magnets before about 2014.

Figure 2: Projected time-scale for an LHC upgrade.

A new insertion region lay-out is considered to be one straightforward way to raise the LHC luminosity. A new lay-out should allow for a lower $\beta^*$ and a reduction in long-range beam collisions. Various options for a new insertion region are being considered:

- Relatively simple designs include the LHC baseline insertion region with larger bore quadrupoles. This is referred to as the Quadrupole First variant. Moreover, the Dipole First option has the separation dipoles between the interaction point and 100 mm large-aperture quadrupole magnets. This option yields fewer long-range collisions but a larger $\beta_{\text{max}}$. These two designs have the potential of reducing $\beta^*$ by between a factor of 2 to 3.

- Alternative designs include incorporating the low-$\beta$ quadrupoles between separation magnets and lay-outs with large crossing angles. Such designs might reduce $\beta^*$ by between a factor of 2.5 to 5.

Re-designing the insertion regions offers the possibility of moving elements of the accelerator, such as the low-$\beta$ quadrupoles, separation dipoles and absorbers, closer to the interaction point to gain even more in luminosity. Integration with the experiments and their shielding should be studied.

Since an upgrade to the insertion region cannot alone achieve a factor 10 increase in the luminosity, additional ways must be developed to maximise the luminosity. Table 1 summarises example parameter sets whereby luminosities approaching $10^{35}$ cm$^{-2}$ s$^{-1}$ are reachable.

Additional gains may be obtained by upgrading the LHC injection chain in order to increase the beam intensity and brilliance beyond their present ultimate values. Upgrading the proton LINAC from 50 MeV to 120-160 MeV would reduce space-charge effects at the PS Booster injection, facilitating delivery of the highest LHC intensities. Equipping the LHC with superconducting magnets and upgrading the transfer...
lines to allow injection in to the LHC at 1 TeV would increase the peak LHC luminosity by nearly a factor of 2 and would be the first step needed in view of an LHC energy upgrade. Finally, the SPL or a rapid-cycling synchrotron could be included in the LHC injector chain.

Table 1: Summary of example parameter sets whereby luminosities approaching $10^{35} \text{ cm}^{-2} \text{ s}^{-1}$ are reachable.

<table>
<thead>
<tr>
<th>parameter</th>
<th>symbol</th>
<th>normal</th>
<th>ultimate</th>
<th>shorter bunches</th>
<th>longer bunches</th>
<th>superlunch</th>
</tr>
</thead>
<tbody>
<tr>
<td>bunches</td>
<td>$n_b$</td>
<td>2001</td>
<td>2008</td>
<td>5016</td>
<td>901</td>
<td>1</td>
</tr>
<tr>
<td>bunch spacing</td>
<td>$\lambda$</td>
<td>1.36</td>
<td>1.67</td>
<td>1.7</td>
<td>6.0</td>
<td>6000</td>
</tr>
<tr>
<td>number of bunches</td>
<td>$n_b$</td>
<td>20</td>
<td>22</td>
<td>12.5</td>
<td>71</td>
<td>18000</td>
</tr>
<tr>
<td>norm-transverse</td>
<td>$\varphi_{0}$</td>
<td>1.9</td>
<td>2.6</td>
<td>2.5</td>
<td>1.7</td>
<td>1.0</td>
</tr>
<tr>
<td>$\beta^*$</td>
<td>$\beta^*$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
</tr>
<tr>
<td>$\sigma^*$</td>
<td>$\sigma^*$</td>
<td>Gaussian</td>
<td>Gaussian</td>
<td>Gaussian</td>
<td>uniform</td>
<td>uniform</td>
</tr>
<tr>
<td>$\sigma_z \ [\text{mm}]$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
</tr>
<tr>
<td>$\sigma_x \ [\text{mm}]$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
</tr>
<tr>
<td>$\sigma_z \ [\text{mm}]$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
</tr>
<tr>
<td>$\sigma_x \ [\text{mm}]$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
</tr>
<tr>
<td>$\sigma_z \ [\text{mm}]$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
</tr>
<tr>
<td>$\sigma_x \ [\text{mm}]$</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
<td>3.55</td>
<td>3.75</td>
<td>3.75</td>
</tr>
</tbody>
</table>

TECHNOLOGICAL CHALLENGES FOR AN LHC LUMINOUSITY UPGRADE

The upgrade to the LHC luminosity requires R&D on several fronts, including in the following areas:

- Up to 15 T high-field superconducting cables and magnets. The aim is to produce superconducting cable with a non-Cu $J_c$ of up to 1500 A/mm$^2$ at 15 T and a temperature of either 4.2 K or 1.9 K. This will lead to the construction of 15 T dipole magnets and 12 T (100 mm aperture) quadrupole magnets of accelerator type to be used in the upgraded insertion regions.
- Powerful and sophisticated RF devices for beam manipulations.
- Medium-field fast-pulsed superconducting cables and magnets.
- Accelerator design and integration within existing constraints. This includes studies on energy deposition and radiation hardness of the accelerator equipment.

At present, the only serious candidate to succeed the NbTi superconductor used in the LHC is the intermetallic compound Nb$_3$Sn. A series of record-breaking dipole magnet models, operating in the 10- to 15 T range have been developed at Twente University and at LBL. R&D in this area has also started within the CARE-NED Joint Research Activity, which has the aim of developing large aperture (> 88 mm) high-field (> 15 T) dipole magnet models. However, all constructed magnet models are not yet of accelerator class and further work is thus required.

Achieving the same $\beta^*$ with NbTi requires a quadrupole magnet aperture of 120-130 mm and a length of 8-9 m. This would result in a 30% increase in $\beta_{\text{max}}$ and ~20% more long-range collisions. Moreover, the present NbTi technology is not sufficiently radiation-hard and is more sensitive to beam heating.

Furthermore, BNL, FNAL, LBL and Stanford University are collaborating with CARE within the LARP (US LHC Accelerator Research Program) framework to improve the long-term physics research opportunities of the LHC. The main emphasis of the present studies is on developing large-aperture, high-gradient quadrupoles. High-aspect ratio, large-aperture separation dipoles are also being considered. The aim of the work is to deliver accelerator-ready magnet designs in their prototype form in the years 2010-2012.

THE GSI COMPLEX UPGRADE

The upgrade to the GSI accelerator complex – coined the Facility for Antiproton Ion Research (FAIR) - was presented. The accelerator physics and technology for FAIR pose a number of R&D challenges in the fields of high-gradient, low frequency RF cavities, novel lattice/collimation design, superconducting, fast-ramping synchrotron magnets, ultra high vacuum for intense beams, control of collective effects and fast stochastic and electron cooling. Several common areas of R&D between the LHC upgrade and FAIR were identified.

CONCLUSIONS

An upgrade to the LHC luminosity will provide a significant extension to the LHC physics reach. The challenge for the both the accelerator complex and experiments to reach a luminosity of $10^{35} \text{ cm}^{-2} \text{ s}^{-1}$ is, however, considerable. In view of this, it is recommended that the following points are taken under consideration:

- Several alternative ways to achieve the high luminosity goal must be considered.
- The R&D directions must be chosen judiciously.
- The R&D effort must start now.
- As resources are limited, convergence among the various options is a prerequisite.

ACKNOWLEDGEMENTS

The author would like to thank the organisers of the workshop for the stimulating and interesting forum.
Abstract

We summarize the Session 2 of the HHH-2004 workshop, which was devoted to beam dynamics and intensity challenges. The topics discussed included the extrapolation of collimation system, machine protection, beam dump, and rf systems to various LHC upgrade scenarios, space-charge studies for the GSI complex, electron-cloud effects at RHIC and SNS, and a review of vacuum physics, which addressed, in particular, the possible interplay of ions and electron-cloud effects.

1 INTRODUCTION

The CARE-HHH-APD Workshop on “Beam Dynamics in Future Hadron Colliders and Rapidly Cycling High-Intensity Synchrotrons” (HHH-2004) was held at CERN, November 8–11, 2004. The workshop discussed possibilities to boost the LHC luminosity by up to a factor of 10 after about 7 years of operation, i.e., around 2014. Such LHC upgrade would include new interaction regions, moderate modifications of several subsystems to cope with a higher beam current, and, possibly, either long-range beam-beam compensation or crab cavities. Various distinct upgrade scenarios are presently being considered.

In Session 2 of the HHH-2004 workshop, 7 presentations were given, namely

- Collimation by R. Assmann,
- Machine Protection by K.-H. Mess,
- Beam Dump by B. Goddard,
- Space Charge and Optics Studies for the GSI Complex by G. Franchetti,
- Electron-Cloud Effects by J. Wei,
- Vacuum, Beam Scrubbing, and Pressure Rise by O. Grobner,
- RF & Feedback Systems for Bunch Shortening by J. Tuckmantel.

In the following we describe some highlights from the various presentations, with emphasis on aspects relevant to the upgrade.

2 HANDLING & DISPOSING OF ENERGETIC HIGH-INTENSITY HADRON BEAMS

For the nominal LHC parameters, the stored energy in each beam is about 400 MJ. This is two orders of magnitude higher than in operating hadron storage rings like Tevatron or HERA. R. Assmann showed that the LHC luminosity upgrade implies a further increase in stored beam energy by a factor 2–3, which appears quite moderate compared with the step from Tevatron to the nominal LHC. It is thought that this increase can be handled after having gained experience with collimation during several years of LHC operation. In case of an irregular asynchronous beam dump, the energy impact on the collimators is a concern. Depending on the specific upgrade scenario this energy impact increases either by a moderately 50%, by a factor of 4, or, in the case of superbunches, by 2 orders of magnitude. The latter value does not seem manageable. R. Assmann pointed out that beyond half the nominal intensity, an upgrade of the present phase-0 baseline collimation system will be necessary. This upgrade may involve metallic collimators, high-Z materials, improved precision, or altogether new technologies or alternative schemes, such as crystals, renewable collimators, plasmas, or non-linear collimation. Advanced collimation is studied partly in collaboration with US-LARP and other European partners.

K.-H. Mess discussed whether an upgraded LHC can be protected. He considered three different upgrade scenarios: (1) an intensity upgrade from 0.58 A per beam to 1.7 A, (2) an energy upgrade from 8.33 T to 9 T magnetic field with the existing magnets, and (3) an energy upgrade from 8.33 T to 15 T with new magnets. He stressed that the beam intensity creates dangerous problems already for the nominal beam and that a further intensity or energy increase will exacerbate the requirements on the collimation system. He strongly discouraged the moderate energy increase, in view of the dramatic decrease in temperature margin this would entail. He concluded that an energy increase with stronger magnets is not precluded, but that this would require a significant space in the tunnel. Alternatives for the cold diodes would be needed. A possible option is the Tevatron solution with HTS switches. For quench protection a sector may have to be split into 8 subdivisions, interleaved with dump resistors for energy extraction, rather than into 2 as at present, in order to limit the voltage drop across the magnets. However, as a bottom line, there is no fundamental obstacle from the viewpoint of machine protection.

B. Goddard discussed the implications of an LHC intensity or energy upgrade for the beam-dump system. The latter consists of 4 components: (1) extraction system, (2) dilution system, (3) beam-dump absorber, (4) protection devices. An upgrade of all these systems for higher beam intensity looks feasible. The material of beam dump and protection would have to be modified by reducing the carbon density, and the sweep length of the dilution system...
would need to be doubled. For a beam energy of 14 TeV, the lengths of dump absorber, dilution system, and protection devices would have to be quadrupled, which is challenging. For the superbunch scenario, an adequate upgrade of the dump system seems impossible, especially for the dilution system and the protection devices. With normal bunch structure, a four times larger sweep length would allow for a maximum current of up to 3 A at 7 TeV and 0.8 A at 14 TeV. In particular, operating the beam dump at 1.7 A current looks feasible.

3 ELECTRON CLOUD AND VACUUM

The electron-cloud experience at RHIC was reported by J. Wei. RHIC is a superconducting storage ring, whose performance is currently limited by the pressure increase due to electron cloud. Though the pressure rise is mostly seen in the warm sections, multipacting has also been observed in the cold arc. A number of detectors installed in the warm regions show electron signals. Tune shift and beam instabilities possibly associated with the electron cloud are noticed as well. The onset of the pressure rise in RHIC strongly depends on the bunch spacing. Halving the spacing from 216 ns to 108 ns, reduces the threshold by almost a factor of two. The electron signal measured by a dedicated detector is almost perfectly proportional to the logarithm of the pressure measured in the same region of the ring. J. Wei also discussed expectations and precautions for the accumulator ring of the Spallation Neutron Source. Here electrons are assumed to be generated at the stripping foil, at the collimators by beam loss induced emission, and, by a ‘trailing-edge multipacting’ mechanism at other places around the ring. A large number of preventive measures were adopted, among others surface coating by TiN, beam-in-gap kicker, solenoids for the straight sections, clearing electrodes, e.g., dual-function BPMs, and a tapered magnetic field near the injection foil. If necessary, Landau damping can be enhanced by painting or exciting nonlinear magnets. J. Wei noted that the storage ring has higher beam intensity, but a smaller aperture than an equivalent rapidly cycling synchrotron, and that the small aperture counteracts the trailing-edge multipacting.

O. Grobner discussed a large variety of vacuum aspects for an LHC upgrade. Experiments at EPA have demonstrated that the dependence of the dynamic pressure increase on the beam energy is different for warm and cold sections. At room temperature, the increase is proportional to the 4th power of the beam energy, while in the cold arcs it is expected to depend on the 3rd power. In either case the pressure rise also varies linearly with the beam current. The desorption yield can be suppressed by conventional getter coating, which yields a factor 10 improvement compared with baked stainless steel (\(< 10^5\) molecules/photon). Coatings with special NEG films as developed at CERN provide a further reduction in the desorption yield (\(< 10^6\) molecules/photon) in addition to a high pumping speed. SPS measurements of the electron flux at the wall of a cold detector by V. Baglin showed 50 mA/m current with a positive bias and 3 mA/m with a negative bias. O. Grobner interpreted this latter value as an ion current, which thus amounts to a significant fraction (more than 5%) of the electron current. Using well-known ionization cross sections and the expression for a standard ionization pump, he estimated a surprisingly large electron path length of \(10^6\) m. The pressure rise in the LHC results from a combination of processes, namely thermal outgassing, electron-stimulated desorption, and primary plus secondary ion desorption. The equilibrium gas density as a function of beam current is then obtained as a rather complicated expression, a generalization of the classical pressure-bump equation. O. Grobner pointed out that the secondary ionization by the electron cloud can reduce the critical beam current for the pressure-bump instability, that the secondary ionization due to the electron cloud is pressure dependent, and that beam-induced multipacting can itself trigger a pressure instability. Making a quantitative estimate, based on the ion current measured in the SPS, he found that the secondary ionization term, from electrons, is about 7 orders of magnitude larger than the primary term describing ionization by the beam. In conclusion dynamic vacuum issues should remain high on the priority list, secondary ionization of the residual gas is linked to the electron cloud, and interplay of ions and electrons is one possible explanation of the long-term memory effect. More studies will be needed concerning the link between electron cloud and ion induced instability. O. Grobner recommend to include ion contributions in the heat load budgets and to confirm these by calorimetric measurements.

4 RF UPGRADE FOR SHORT BUNCHES

J. Tuckmantel discussed beam manipulations and changes to the rf system required for the baseline LHC upgrade scenario with half the nominal bunch length. Shorter bunches allow recovering the geometric luminosity loss, that would otherwise be encountered for a smaller \(\beta^*\) and a constant crossing angle. A brute force approach would simply increase the rf voltage. Since the dependence in the 4th power, the number of cavities would need to be increased 16-fold to 128 units in total. This would have severe cost and space implications and, hence, it is not an attractive option. A more elegant approach is to add an rf system at a higher frequency. For various reasons, a frequency of 1.2 GHz is chosen, which is 3 times the nominal frequency. A voltage of 43 MV is needed for the 1.2 GHz system. However, the bucket area is reduced in this scheme. Therefore, the longitudinal emittance must be limited to 1.77 eVs instead of the nominal 2.5 eVs. For the nominal LHC, at injection the emittance is only 1 eVs. It is later blown up during acceleration in order to stabilize the beam. The upgrade scenario considered implies that the magnitude of the emittance blow-up be reduced. Even with the smaller emittance, the bunch occupies a larger fraction of the (smaller) rf bucket than in the nominal LHC. High-field
1.2–1.3 GHz cavities can be manufactured from bulk Nb, using TESLA technology. The cavity shape may essentially be scaled from the 400-MHz version, except for bolts and wall thickness, which do not scale trivially and would need to be optimized. For the existing 400-MHz 4-cavity module, the counterpropagating beam is passed outside of the cavities and the beam pipe of the accelerated beam. For the 1200-MHz module a set of 12-cavities for each beam could be mounted side-by-side in the same cryostat, offering a space-saving and economical solution. Other layouts are also conceivable. J. Tuckmantel summarized that the technology is ready and reliable, about 15–20 m tunnel length would be needed, and no fundamental problem was found. There are a few technical challenges, all of which appear solvable, like a multitude of outlets, a ‘crowded’ rf system, allocation of power converters and wave guide passage into the tunnel. Beam stability issues and intrabeam scattering with the reduced longitudinal emittance remain to be evaluated.

5 SPACE CHARGE AND OPTICS STUDIES FOR THE GSI COMPLEX

G. Franchetti underlined that the GSI FAIR project is entering a new regime of beam dynamics. Historically there was a distinction between high-energy machines, storing beam for more than $10^6$ turns, in which dynamic aperture and resonances were major concerns, and high-intensity machines, storing a beam for about $10^4$ turns only, where space charge was the dominant effect. The new SIS100 ring will operate at high intensity over $10^5$ turns. The fundamental questions to be addressed are whether the space charge in this regime leads to new effects and whether a design of a high intensity ring can be based on single-particle dynamics. An ongoing collaboration between CERN and GSI performs experiments in the CERN PS, aimed at benchmarking the related computer simulations. The linear and nonlinear optics of the PS is taken into account [1]. Both experiments and simulations reveal two distinct regimes, separated in tune space, where either the emittance blows up or a part of the beam is lost. The observations and measurements agree well for the emittance dominated regime and within a factor 2–3 for the loss prediction. In the beam-loss dominated regime, the longitudinal shape of the bunch changes from a Gaussian to one with triangular edges. The interpretation, supported by computer simulations, is that particle at large longitudinal amplitudes are crossing the 4th order resonance during the synchrotron motion. These particles can be trapped on the resonance and lost. Similar experiments were also conducted at the existing GSI SIS18 ring. These experiments also observed particle loss and bunch-length shortening, which is explained by the same trapping and loss mechanism. In conclusion, G. Franchetti and co-workers developed a code for describing long-term effects in a high-intensity bunched beam with a nonlinear lattice. Space-charge induced emittance growth (halo) and long-term particle loss by resonance trapping were studied in simulations and experiments. The consequence for the working point choice is that there exist two regimes which should be avoided. Both of these are located above the resonance in the tune diagram. The closer regime is the beam-loss regime. Adjacent to this and further away from the resonance, one finds a regime with large emittance growth. The working point can be placed on either side of the resonance outside of these two forbidden areas.

6 CONCLUSIONS

No principal showstoppers were found for any of the more conventional LHC upgrade options, in particular for a higher beam current. The only prerequisite of an LHC intensity upgrade is that the collimation system be upgraded too, e.g., by means of some advanced concept. Indeed the properties and the efficiency of the collimation system will prove crucial for any upgrade (as already the case for reaching the nominal LHC parameters). The extrapolation from nominal to upgraded LHC in terms of beam power or energy density appears quite moderate compared with the extrapolation from the Tevatron to the nominal LHC. The superbunch option is ruled out from the accelerator point-of-view, considering the excessive and unsolvable requirements which such bunches would place on the beam dump system and on the collimator survival in case of an asynchronous dump-kicker firing. Electron cloud is an effect which presently limits the performance of RHIC. For the SNS an impressive suite of countermeasures were implemented. The interplay of ions and electron cloud may be larger than previously estimated. Dynamic vacuum pressure and the combination of beam-induced multipacting, electron cloud, and pressure instability should remain high on the priority list and certainly may become more important still for an LHC upgrade. An uncritical item is the rf system needed to shorten the bunch length by a factor of 2, for the baseline upgrade scenario, where a technical solution exists.

Lastly, space-charge effects occurring over long time periods push the frontier of beam dynamics at the GSI FAIR Project. In joint CERN-GSI experiments, pertinent simulations were successfully benchmarked against experiments, identifying two forbidden areas in tune space.

7 ACKNOWLEDGEMENT

We thank all the speakers of Session 2 for their informative, well prepared and clear presentations.

8 REFERENCES


SUMMARY: LHC IR UPGRADE AND BEAM CHOICES

S. Peggs, BNL, O. Brüning, CERN

INTRODUCTION

The spectrum of potential upgrade scenarios shown in Table 1 holds the possibility of boosting the LHC luminosity by as much as a factor 10 beyond the nominal value of $1.0 \times 10^{34} \text{cm}^{-2} \text{s}^{-1}$. All scenarios incorporate new IR magnets that go beyond the present state-of-the-art, such as stronger or larger-aperture low-beta quadrupoles, or specialized beam separation dipoles. Some of them also have significant implications for other hardware sub-systems, such as beam-beam compensators, crab cavities, acceleration cavities, collimators, cryogenics, and beam dumps. Significant upgrade requirements may even extend back up the injector chain.

The presentations made in Session 3 of the workshop investigated these implications in some depth. This summary merely introduces the major themes and issues, leaving the reader to find the details in the Session 3 papers contributed to these proceedings [1, 2, 3, 4, 5, 6].

<table>
<thead>
<tr>
<th>scenario</th>
<th>$L$</th>
<th>$\Delta T$</th>
<th>$N_i$</th>
</tr>
</thead>
<tbody>
<tr>
<td>nominal</td>
<td>1.0</td>
<td>25</td>
<td>23</td>
</tr>
<tr>
<td>ultimate</td>
<td>2.3</td>
<td>25</td>
<td>54</td>
</tr>
<tr>
<td>IR upgrade</td>
<td>4.6</td>
<td>25</td>
<td>108</td>
</tr>
<tr>
<td>super-bunch</td>
<td>9.0</td>
<td>$9 \times 10^4$</td>
<td>$9 \times 10^5$</td>
</tr>
</tbody>
</table>

Table 1: Luminosity $L$, bunch spacing $\Delta T$, and the number of events per bunch crossing $N_i$, tabulated for a selection of different LHC scenarios [6].

REPRESENTATIVE UPGRADE LAYOUTS

At least 5 plausible upgrade layouts are currently under discussion [1]. Three of the most representative layouts are shown in Figure 1.

The “quadrupole first” layout (with a small crossing angle) is shown at the top of Figure 1. Both beams go through a single bore of a quadrupole triplet in this layout, which is the nominal implementation of the LHC as it will first be run. Stronger quadrupoles would allow the focusing center to move closer to the interaction point (IP), allowing a smaller $\beta^*$ and more luminosity for a given number of bunches, and bunch intensities, et cetera. Or, larger bore quadrupoles with the same gradient would hold the slot lengths constant, allowing a quadrupole-by-quadrupole progressive upgrade. In either case the beams will suffer a relatively large number of parasitic long range beam-beam interactions, until the beam reach the first beam separation dipole, approximately 60 meters from the IP.

The “dipole first” layout shown in the middle of Figure 1 separates the beams into separate triplet quadrupole...
bores. This scheme has the advantage of eliminating the long range beam-beam collisions beyond about 23 meters from the IP, but it has the optical disadvantage of moving the triplet focusing center farther from the IP. The question of where and how to absorb the many kilowatts of luminosity debris power is a significant challenge in all scenarios [1]. One suggestion is to use magnetized absorbers in the middle of the IR optics. The dipole first layout has the simultaneous advantage and disadvantage of not needing magnetic absorbers, since the first dipole will absorb much of the debris power.

Finally, the “large crossing angle” scenario shown at the bottom of Figure 1 has the dual advantages of almost completely eliminating long range beam-beam interactions, while maintaining a relatively close triplet focusing center. However, it has the disadvantage of requiring a large number of crab cavities. This in turn stresses the need to reduce the transverse size of the first focusing quadrupoles, in order to reduce the crossing angle (and unburden the crab cavity system) as far as possible.

### BEAM-BEAM COMPENSATION

The nominal bunch spacing of $\Delta T = 25$ ns leads to one long range beam-beam interaction every 3.75 m up to a distance of $L_{sep}$ from the IP, when the first beam separation dipole is encountered. Thus there are as many as 30 interactions per IR [2]. Each IR generates a long range beam-beam tune shift of

$$|\Delta Q| \approx \frac{2L_{sep} \xi}{c \Delta T} \left(\frac{\theta}{\sigma^*}\right)^2$$

(1)

where $\xi$ is the (head-on) beam-beam parameter, $\theta$ is the total crossing angle, and $\sigma^*$ is the RMS angular size of the beam at the IP. The plane of the crossing can be alternated between the 4 IRs in the LHC, leading to a significant cancellation of the net tune shift. This cancellation even works for PACMAN bunches, near bunch pattern gaps, that experience non-standard beam-beam collision sequences.

### CRAB CAVITIES

When the crossing angle is larger than the natural aspect ratio of the bunch, $\theta \geq (\sigma^*/\sigma_s) \sim 1$ mrad, then a lot of luminosity is lost because the head of one bunch does not collide with the tail of the other. Figure 4 illustrates the principle by which transverse deflecting mode RF crab cavities fix this problem, inducing a localized perturbation in
Figure 4: The crab crossing principle. Incoming bunches are tilted by transverse deflecting mode crab cavities on the extremities of the IR so that they collide head-on. The tilt is removed on exit by another set of RF cavities [2].

the closed orbit as a function of longitudinal displacement from the center of the bunch.

A single crab cavity will begin engineering tests in one of the KEK B factory rings in late 2005, followed by an operational test with one cavity in each ring in 2006 [3]. Table 2 compares KEK B crab cavity parameters with sample LHC parameters, showing that the LHC implementation is much more extreme. The total RF voltage required is given by

\[ V = \frac{cE \tan(\theta/2)}{2\pi e f_{RF} \sqrt{\beta^* \beta_{crab}}} \]  

(2)

If the total crossing angle is \( \theta = 8 \) mrad, then the beams are only separated by about 18 cm at a distance of 23 m from the IP, the closest approach of the first (side-by-side) quadrupoles.

Crab cavity phase errors generate transverse displacements at the IP which, when coupled with the beam-beam effect, can lead to unacceptably large emittance growth rates. This effect is unimportant in electron colliders like KEK B, where it is suppressed by synchrotron radiation damping. Figure 5 shows emittance and luminosity evolution from a short timescale simulation, also including simulations of acceleration cavity phase noise for comparison [3]. It remains an open issue whether acceptably small emittance growth rates are possible. For this reason it has been suggested that, if the 2006 KEK B crab cavity tests are successful, then a crab cavity should be installed and tested in a hadron machine to demonstrate a level of RF phase noise compatible with acceptable emittance preservation.

**LONGITUDDINALLY FLAT BUNCHES**

It is predicted that, under a set of conditions and assumptions including a large crossing angle \( \theta \) and operation at the beam-beam limit, the luminosity with longitudinally flat beams is \( \sqrt{2} \) larger than with longitudinally gaussian beams [2, 5]. Figure 6 illustrates the extreme case, in which a single flat super-bunch extends around almost all of the circumference, contained (and accelerated) by barrier bucket induction cavities. Induction acceleration has recently been demonstrated at the KEK proton synchrotron [4].

A single full length super-bunch is not optimal. Figure 7 predicts that under some conditions the optimal total bunch length is about 280 m, divided into a convenient number of bunches. Some suggest that the technology for confin-

---

**Table 2: Comparison of KEK B crab cavity parameters with those typically required for an LHC upgrade [2, 3].**

<table>
<thead>
<tr>
<th></th>
<th>KEK B</th>
<th>LHC</th>
</tr>
</thead>
<tbody>
<tr>
<td>crossing angle</td>
<td>( \theta ) [mrad]</td>
<td>22</td>
</tr>
<tr>
<td>beam energy</td>
<td>( E ) [TeV]</td>
<td>0.008</td>
</tr>
<tr>
<td>collision beta</td>
<td>( \beta^* ) [m]</td>
<td>0.33</td>
</tr>
<tr>
<td>crab beta</td>
<td>( \beta_{crab} ) [km]</td>
<td>0.1</td>
</tr>
<tr>
<td>RF frequency</td>
<td>( f_{RF} ) [GHz]</td>
<td>0.51</td>
</tr>
<tr>
<td>RF voltage</td>
<td>( V ) [MV]</td>
<td>1.4</td>
</tr>
</tbody>
</table>
Figure 6: Bunch confinement and acceleration using conventional resonant RF and induction cavity RF systems. LEFT: many bunches in a conventional system. RIGHT: a single long super-bunch under the influence of barrier pulse voltages generated by an induction system [4].

ing “intermediate length” bunches using a conventional RF system operating at harmonics of 40 MHz is more appropriate than an induction RF system [5]. A larger number of shorter bunches is also favored by the experimentalists, who are daunted by the extraordinary event pile-up difficulties that arise for super-bunches (see the last row in Table 1). Figure 8 shows that sufficiently flat beams are produced using only a 3 harmonic system, whatever the fundamental frequency [5].

Figure 7: Luminosity in the LHC versus total beam length $n_b l_b$ and line density $\lambda$. The vertical surfaces represent constraints imposed by a maximum total beam-beam tune spread of 0.01 (flat surface) and a maximum average current of 1.1 A (curved surface) [5].

Figure 8: Longitudinal phase space and line density of a bunch held by three RF systems, for example 40, 80, and 120 MHz. The distribution is flat enough to reap the predicted advantages of longitudinally flat beams [5].

use NbTi conductors (arguably more appropriate in the near term), or Nb$_3$Sn conductors (in the medium term). In any case, some years of initial operation of the LHC will be required before settling on the optimum upgrade scheme. This does not obviate the need to pursue further studies and R&D now, so as to be ready then. It is also possible that there will be multiple (perhaps modest) upgrades in sequence, and not one single monolithic upgrade.

ACKNOWLEDGMENTS

We are very grateful for the help that the participants in Session 3 gave to us in the preparation of this summary.

REFERENCES


CONCLUSIONS

There is much activity exploring a spectrum of potential LHC upgrades that have the potential to enhance the luminosity by as much as an order of magnitude. Many technologies in different arenas need investigation. Nonetheless, all scenarios incorporate new IR magnets that go beyond the present state-of-the-art, whether these magnets
HHH-2004: SUMMARY OF SESSION FIVE*

W. Scandale#, CERN, Geneva, Switzerland

Abstract

In the CARE–HHH-APD workshop, session 5 was devoted to ‘Fast cycling injector’. The scope was given by the more and more compelling request to investigate possible upgrade of the injector chain for existing or future hadron accelerators in Europe. The main goal was to review specific problems related to the upgrade of the LHC injector chain and to the approved FAIR project at GSI. Six talks were devoted to investigate hot problems in this field.

SUMMARY OF THE SESSION

In the CARE workshop ‘HHH-2004’, session 5 hosted six talks reviewing some of the critical problems related to the FAIR project at GSI and to the upgrade of the LHC injector chain.

In the first talk, Roland Garoby from CERN analysed the needs and the means for reaching higher brightness beams in the LHC injector chain. He pointed out three possible ways to reach this goal: (i) the use of batch compression mechanisms in the CERN-PS, based on a sophisticated RF gymnastics, can provide up to 2.6 $10^{11}$ protons per bunch within the nominal emittance, thereby increasing by more than 50 % the beam brightness at the PS ejection; (ii) the use of LINAC 4, providing 160 MeV/c proton beams to the CERN-PSB, together with some RF gymnastics in the PS, can provide up to 3.0 $10^{11}$ protons per bunch within the nominal emittance and almost double the beam brightness at the PS ejection; (iii) the construction of the SPL, replacing the CERN-PSB, can provide a further increase of the performance, providing up to 4.0 $10^{11}$ protons per bunch within the nominal emittance at the PS ejection. Roland also mentioned the possibility of replacing the CERN-PSB by an appropriate Rapid Cycling Synchrotron (RCS), having the same performance as the SPL and most likely a more favourable cost-effectiveness.

In the second presentation, Peter Spiller from GSI presented the status of the two-stage synchrotron project SIS 100/SIS 300. SIS 100 is the high intensity compressor stage whilst SIS 300 is the high-energy stretcher stage of the accelerator complex for heavy ions and antiprotons planned during this decade at GSI. The two rings will use very advanced pulsed superconducting magnets still requiring validation studies, implying large international collaborations. In addition, an intense R & D program on the RF design and components should provide the expected performance for bunch compression. Other issues under investigation are the vacuum system, the lattice optimization and the collimation efficiency.

The third talk, given by Gebhart Moritz from GSI, presented the status of the superconducting magnet program for FAIR. Very ambitious parameters are requested both in SIS 100 (ramp of 4 T/s) and in SIS 300 (top-field of 6 T and ramp of 1 T/s), with severe implications on eddy current and persistent current effects, on cryogenic load and stabilization, on mechanical structure and expected lifetime of magnets, on quench protection and on radiation hardness. A superconducting nucleotron type dipole is proposed for SIS 100, although a superferric variant is still not excluded. Gebhart also presented the results of the GSI01 dipole build at NBNL, based on the RHIC cable and magnet design. In this magnet, the observed and the computed thermal loss do not fully coincide and the final validation of electro-magnetic and cryogenic loss model is yet under scrutiny. The SIS 300 dipole is presently based on a conceptual design study made at IHEP in Protvino, derived from the former UNK dipole. The SC cable and the SC filament are still under investigations, in particular to decide the filament size, the parameters of the embedding copper matrix and the shape and the characteristics of the final cable.

In the fourth presentation Elena Shaposhnikova from CERN recalled that impedance is the present cause of intensity limitation in the CERN-SPS. In 1999/2000 RF cavities for leptons and other components no more in use were removed, whilst pumping ports and kickers still in use were shielded, thereby increasing the SPS wall smoothness. The effect on longitudinal impedance was observed in 2001 through a reduced bunch spectrum response, a reduced dependence of the bunch lengthening on the bunch intensity, a reduced quadrupole oscillation frequency and coherent tune shift as a function of the bunch intensity. Present intensity limitations are induced by electron-cloud phenomena, intensity dependent capture loss at injection, residual coupled bunch instability during acceleration and severe requirements on bunch shaping at flat-top, imposed by the LHC requirements. Additional investigations will be required to fully master the residual beam loss in the SPS.

In the fifth presentation, Massimo Giovanniuzzi from CERN presented the idea of injecting and extracting particles with reduced loss through stable resonant island adiabatically moving in the phase space. Massimo proposed to use the novel method for extraction/injection in the beam transfer from the PS to the SPS, where at present loss of up to 15 % are produced at the PS electrostatic septum. The method consists in creating and populating stable resonant islands, which can adiabatically bring particles at large action by a tune change. Experimental results in 2002 and 2003 showed the effectiveness of the methods and a considerable reduction of particle loss during extraction. A similar mechanism can be used for injection process as well. On going studies will clarify how one can better control the core and the island beam population, in order to flatten the beam density during the transfer process.

In the last talk, Frédéric Méot from CEA/DAPNIA presented an overview of Fixed Field Alternate Gradient
Synchrotrons. From the old-days of the Mark II project at MURA to the present days of the scaling/non-scaling linear/non-linear design studies, several basic improvements have been proposed to the FFAG concept. In the spiral sector version, a positive field superimposed to the alternating radial sector field shall reduce the circumference factor. At KEK, the POP proton machine, based on high field non-linear gradient is intended as a proof of principle, whilst a 150 MeV proton synchrotron is already operational since 2003. At FNAL, a 0.5 MW 60 MeV FFAG with radial sectors is proposed as a 105 Hz proton driver. At KEK, a spiral sector version at 750 Hz should produce 1.5 GeV 10 MW beams. The Japanese neutrino factory proposes a variant with a 50 GeV FFAG ring. Finally at NBNL, a 1.5 GeV FFAG is proposed as a pre-injector for the AGS to improve the intensity performances. In conclusion the FFAG concept is considered cost effective both for multi-GEV multi-MW fast cycling proton drivers and for compact low-energy accelerator, i.e. for medical applications.

CONCLUSIVE REMARKS

The presentations and the subsequent discussions pointed out in a clear manner that to enlarge the reach both in nuclear and high-energy physics, high intensity/high brightness proton beam are largely required at several energy scales. The open points requiring additional R & D cover a vast range of technologies, from superconducting wires, cables and magnets producing rapid cycling fields to RF components providing large accelerating fields with a reduced impedance budget, from dynamic vacuum to beam dynamics control. Mandatory requirements are those of preserving reduced impedance budget and small beam emittance and controlling beam loss not only with robust collimators but also introducing ‘loss-free’ injection/extraction mechanisms. The tasks are heavy and should be shared in the nuclear and high-energy physics communities. Networking activities and cooperative approaches to R & D will be very effective ways to face the new challenges.
Panel Discussion on Single Particle Codes

O. Brüning, CERN, Geneva, Switzerland

Abstract

The panel discussion on single particle codes included six panel members (O. Brüning [CERN], W. Decking [DESY], F. Meot [CEA], K. Ohmi [KEK], F. Schmidt [CERN], J. Wenninger [CERN]) who contacted 12 experts prior to the discussion at the HHH workshop (R. Assmann [CERN], H. Groth [CERN], W. Herr [CERN], E. Keil [CERN], S. Peggs [BNL], F. Plat [BNL], G. Robert-Demolaize [CERN], S. Redaelli [CERN], F. Ruggiero [CERN], R. Talman [CORNELL], T. Sen [FERMILAB], F. Zimmermann [CERN]). The panel discussion included 2 presentations, one submitted presentation that was not shown during the discussion, 2 email responses from the contacted experts and input from other session of this HHH workshop. Since the panel discussion included only a small group of selected people with only limited input from the general audience it does not reflect necessarily the accelerator community as a whole but rather the consensus of a discussion within a selected group of experts.

QUESTIONS

The panel members identified 6 main topics and goals for the the panel discussion:

- Generate a catalog of existing codes.
- Address the problem how one can achieve transparency between the input and output formats of the different Single Particle Codes and discuss the status of SXF?
- Discuss the options for combining single particle tracking with general tools for non-linear dynamics: what compromises in speed, flexibility, versatility and maintainability are still practical?
- What are the possibilities for real time simulations and what compromises between speed and accuracy are acceptable?
- What is the 'wish list' for additional features in Single Particle Tracking tools in the accelerator community (e.g. tracking with dynamic effects of the magnet field errors, tracking with acceleration, tracking with aperture limitations and collimators and tracking with collective effects such as space charge)?
- Discussion on the managerial issues for the code development and maintenance in large inter laboratory collaborations.

THE MAIN PRESENTATIONS

Overview of existing Single Particle Simulation Programs

W. Decking defined Single Particle Codes in his presentation as simulation programs that describe the motion of a particle in the 6 dimensional phase space under the influence of external fields [1]. This definition was used throughout the panel discussion and W. Decking provided a list of 26 codes currently used in the accelerator community. Furthermore, W. Decking put forward an idea by N. Waker from DESY, to divide any accelerator project into three main phases:

- Basic Lattice Design: this phase relies mainly on simple models (sequence of elements) with generic magnet families and classes and requires fitting routines for matching the optics parameters of the machine.
- Performance Simulations: this phase requires more complex models that allow the simulation of machine imperfections and a refinement of the key parameters using tuning algorithms and diagnostic tools.
- Modeling of the Real Machine: this phase requires fast simulation modules and diagnostic tools similar to those present in the machine control room.

Special Requirements for Fixed Field Alternate Gradient Machines

The presentation by F. Meot addressed the special requirements for Fixed Field Alternate Gradient (HFAG) machines where the simulations must be accurate over a wide range of the particle momentum [2]. Because most Single Particle Codes are based on simple transport maps (linear and non-linear) which are optimized for a specific design momentum they can not easily be applied to the design of HFAG machines. The "Zgoubi" method discussed by E. Metral is based on a direct evaluation of the Lorentz equation which provides accurate results for a wide range of particle momenta but requires an accurate magnet field maps for each magnetic element.

MADX as an Example for an Inter laboratory Collaboration

F. Schmidt presented an overview and status report of the MADX project [3]. The MADX project is based on the old MAD8 program and features a truly modular program structure. Its core program is written in C and most modules are written in C, FORTRAN90 and FORTRAN77.
However, other programming languages are in principle acceptable for new program modules (e.g. C++).

CONCLUSIONS AND SUMMARY OF THE MAIN DISCUSSION POINTS

The two presentations by W. Decking and F. Meot underlined the need for specialized codes that are optimized for the special needs and the different phases of a given project. The discussion used the different project phases of the LHC as an example for the special needs in each phase. The discussions underlined the need to combine 'simple' single particle tracking routines with other effects such as rest gas ionization, space charge and impedance effects and the modeling of the machine aperture and the scattering and absorption processes for particles that reach the machine aperture.

The presentation by P. Spiller on the SIS 100 storage ring design for GSI [4] provides a very interesting example for the special needs of an accelerator project. For the SIS 100 design the Single Particle Simulation were combined with a detailed aperture model and an accurate simulation of the scattering and absorption processes for particles that hit the mechanical aperture of the storage ring. This modeling and the search for the highest possible collimation efficiency was the main optimization criteria for the storage ring lattice design.

In order to avoid the use of several different specialized codes it and a 're-invention' of a Single Particle Simulation Program for each project it is necessary to develop general toolkits and one general purpose program with a modular program structure. MAD9 and the CLASSIC collaboration and the new MADX project presented by F. Schmidt present state of the art attempts for such a modular approach. The MADX project is based on the old MAD8 program with a new core module written in 'c' and a truly modular reorganization of the old MAD8 functional units. The MADX program combines general optics and tracking routines with sophisticated packages such as Differential Algebra and post-processing toolkits. Its modular program structure allows straightforward integration of 'new' modules.

However, in spite of its success, the presentation by F. Schmidt showed that the maintenance of the individual MADX program modules by different 'module keepers' that come from different laboratories is still a challenging task and requires special organizational and administrative work. The MADX project has for this purpose a program custodian who looks after the integration of new modules and is ultimately responsible for the release of new program versions.

The discussions also identified the need for transparency of the input and output formats of different existing Single Particle Tracking tools. A universal program language, SXF, was proposed for this purpose some years ago. However, while SXF is used in BNL (not for code comparison but within the RHIC control system) and several programs, such as MADX, support the language, it was not clear how many programs support it and how it should be developed in the future. R. Talman proposed in an email contribution to improve SXF with XML and ADX.

The discussion of a 'wish list' for future program developments underlined that there will be as many new wishes for functionality as there are new accelerator projects. For example the loss simulations with an aperture model and its application to the optics and lattice design for the SIS 100 storage ring is a good example for this aspect. However, the discussion showed also at least three concrete examples for desired future developments:

- The modeling of dynamic effects is clearly a desired option for the design and operation of any super conducting storage ring.
- Simulations including acceleration is an important feature for a wide range of future projects such as linear collider and neutrino factories.
- The incorporation of collective effects is a very interesting option for the design and operation of machines with high beam intensities.

The possibility of real time simulations for online applications in the control room was also identified as an important future development. However, the discussions could not answer what compromises between simulation speed and accuracy are acceptable and what functionality (single particle only or the modeling of some collective behavior) is required. Both points need to be addressed in future discussions concerning the options for a online modeling of the beam dynamics.

REFERENCES

[2] F. Meot in these proceedings.
Summary of panel discussion on beam-beam codes and simulations

W. Herr, CERN, Geneva, Switzerland

Panel members: M. Furman (LBL), K. Ohmi (KEK), T. Sen (FNAL), F. Zimmermann (CERN)

Abstract

This is a summary of a panel discussion following a presentation on beam-beam simulation codes [1]. It reflects the main issues raised and considered important during the discussion by the panel members and the audience. Due to the limited time available for the formal discussion, some relevant deliberations during informal discussions mainly between and with panel members have been included as well.

BEAM-BEAM SIMULATIONS

Beam-beam simulations are an essential tool to design and run a colliding beams facility. They are used in all phases of the design and during the operation, basically until the machines are switched off.

Applications of beam-beam simulations

In the very early phase they are usually used to define the basic performance parameters. In this stage some of the studies are of more “academic” nature and possible problems should also be unraveled and possible solutions proposed. Unfortunately, and this is true for most other simulations, it is rather difficult to predict the future or the unexpected. Surprising effects and the lack of proper simulation techniques and resources have led to the reputation that beam-beam simulations can only predict the past, as one of the panel members has formulated a provocative statement. Recent progress in understanding and advanced techniques has largely improved the predictive power and the limitations are well understood.

In the second stage of the design the use of beam-beam simulations is vital for the optimization of the parameters and the definition of the operational conditions.

During the (usually many) years of operation of a collider the beam-beam simulations become one of the most important tools to understand and improve the performance. This was demonstrated in many machines, e.g. SPS proton-antiproton collider, LEP and the Tevatron. Many observations were only understood with the help of such simulations and helped to improve the simulation programs themselves [1]. Including the much improved computing resources we have now a number of reliable simulations codes and tools available or under development [1].

Types of beam-beam simulations

The beam-beam effects act on single particles as well as on an ensemble and we distinguish two basic types of codes. Weak-strong simulations which are used to study single particle behaviour under the influence of a static beam-beam force and strong-strong simulations where both beams are allowed to change their parameters under the influence of the opposing beam. In the latter case we look for self-consistent behaviour to study collective motion and emittance growth. Both types of simulations are used to evaluate the performance in all phases of a collider.

WEAK-STRONG SIMULATIONS

Typically single particles are subjected to beam-beam forces in this types of simulations.

Single particle simulations

A straightforward and immediate application is the evaluation of the amplitude dependent detuning caused by the opposing beam. The results are usually presented in the form of "tune footprints" and are used to estimated the required tune space in the working diagrams. They also serve as a benchmark for the overall strength of the incoherent beam-beam effects.

The stability of particles under the influence of the non-linear fields is usually estimated by tracking particles through the machine elements which can include the beam-beam interaction as a source of very non-linear fields. The evaluation of the particle loss, dynamic and diffusive aperture and chaotic behaviour allows to estimate the expected stability region.

The definition of tunes, i.e. the working point, of the machine must be done taking the beam-beam effects into account. Since such simulations are often done for machines with imperfections (linear and non-linear, self-consistent optical functions) their correction and the tracking can be integrated into general purpose optics programs.

Fluctuations bunch by bunch

In colliders with many bunches fluctuations must be expected from bunch to bunch and properly taken into account in simulations. Different collision pattern of the bunches can lead to significantly different dynamic behaviour, e.g. due to unsymmetric filling schemes. The properties of ALL bunches must be studied in these cases.

Bunch by bunch effects and self-consistent treatment

Partially as a consequence of bunch to bunch fluctuations but also due to beam-beam effects in colliders with many bunches with different collision schemes, one has to expect different parameters of the bunches. In a weak-strong configuration these parameters can be evaluated us-
ing perturbative methods. However, in a strong-strong situation a perturbative treatment can lead to completely wrong results. In the case of significantly different interaction pattern of many bunches in a collider, the parameters might have to be computed in a self-consistent form when the beam-beam interaction is very strong. Such a self-consistent treatment was essential to understand the LEP performance and the predictions were experimentally verified.

It has been demonstrated that a self-consistent treatment is possible, even for a machine with as many bunches as in the LHC. The computation of parameters such as tunes, orbits and chromaticities for all bunches is reliable even in the presence of intensity or beam size fluctuations and allows to set tolerances on these fluctuations.

**Long term tracking**

The evaluation of slow particle losses or diffusion is in general done using long term tracking, i.e. following the particles over $10^5$ or more turns. To follow a particle for a LHC damping time requires up to $10^6$ turns. Even with the most powerful machines this becomes unrealistic, in particular when many bunches have to be tracked simultaneously. Medium term tracking in the presence of beam-beam effects is however feasible and highly desirable. This can show the effects of beam manipulations carried out over a restricted period of time. Typical examples are ramping of the beam energy, orbit movements, mismatch during injection processes etc. In all these cases the beam-beam effects can play an important role and must be understood qualitatively and quantitatively.

**STRONG-STRONG SIMULATIONS**

Under the name of strong-strong simulations we consider those where both beams are allowed to mutually change each other.

**Coherent effects**

A classical example of a strong-strong simulation is the evaluation of coherent beam-beam modes using multiparticle representing a bunch or beam. Collective effects such as Landau damping or decoherence are automatically included in such programs.

In the case of many bunches like in B-factories or the LHC, the bunches can couple together via (long range and head-on) beam-beam effects and result in multi-bunch modes. Such modes can largely obscure tune measurements or feedback systems and must be controlled or kept small by a proper choice of the configuration and parameters. Rigid bunch models are a powerful tool to study a large range of multi-bunch effects. Since all bunches interact it is a true strong-strong simulation and delivers self-consistent results.

**Emittance growth**

Another important application of strong-strong, i.e. self-consistent simulation techniques is the evaluation of emittance growth due to the beam-beam interaction, possibly in combination with other effects such as machine imperfections, electron cloud, space charge or machine impedance. Such techniques have been used since many years to compute and predict the beam-beam limit in lepton colliders with mixed success. Although the general behaviour and sensitivity to various parameters can be reproduced, a reliable quantitative prediction of the beam-beam limit remains a challenge. In particular in the cases of large colliders or machines with fast radiation damping the results are extremely sensitive to imperfections or initial conditions. However, during the operation of a collider these studies are vital to understand and improve the performance.

**Beam tails and life time**

The evaluation of beam tails and the prediction of the beam life time is one of the most challenging tasks for a simulation program. Conceptual difficulties such as the (relative) small number of particles involved, sensitivity to uncontrollable parameters (initial conditions, imperfections etc.), dependence on history and other effects, make a precise quantitative prediction almost unrealistic. A qualitatively correct description of the life time (or tails) and their dependence on operational manipulations or their responses to parameter changes can be considered a very good success [1].

**6D SIMULATIONS**

The need for the simulation of longitudinal motion is rather obvious for lepton machines where the synchrotron motion is fast and the damping times are short. For most simulation programs applied to hadron machine this is often ignored. Although this may be justified in many cases, some studies might require a full simulation of longitudinal effects. In the case of beam-beam interactions it should be considered whether or not it is required under certain conditions, such as e.g.:

- Presence of a crossing angle
- Offset collisions
- Additional source for Landau damping of coherent modes

**COMPUTATIONAL AND NUMERICAL PROBLEMS**

**Computing resources**

Already single particle tracking for a large number of turns becomes a challenge. The study of coherent effects and emittance growth with multi-particle simulation
programs imposes further difficulties, in particular when multi-particle and multi-bunch simulations are required.

The only solution is the massive use of parallel processing. Several strong-strong simulation programs have already been re-written or modified to take advantage of this possibility (e.g. BEAMBEAM2D, BEAMX, for a list and details see [1, 3]). Single particle tracking codes take advantage of parallel processing by tracking particles simultaneously (e.g. BBSIM, [1]). Depending on the type of computation (e.g. parallel tracking of particles (bunches) or parallelized field solvers) the data transfer between processors can be the ultimate limitation. Newly developed standards like CERN’s OPENLAB with InfiniBand are very promising and are under study for LHC simulations.

Numerical problems

A well-known property of all types of simulations is that not all observed effects have a physical origin but are the result of the numerical treatment of the problem. Unfortunately, there are numerous effects which can make the results irreproducible, imprecise or wrong (i.e. useless). A short and (highly) incomplete list includes:

- Simulation hardware, e.g. different processors
- Simulation software, e.g. different compilers
- Simulation techniques, e.g. different algorithms

To illustrate the last item, in the case of multi-particle strong-strong beam-beam simulations this may include:

- Ill-conditioned problem
- Problem with symplecticity
- Number of macroparticles
- Grid size for field calculations
- Step size for integration etc.
- Choice of the method for field calculations (depends on application: precision and stability ?)
- Validity of approximations (e.g. Gaussian approximation for beam-beam forces [2])

A similar list can be established for weak-strong simulations. Some of these items inject numerical noise into the calculations while others introduce conceptual limitations.

In the problems we consider, the issue of well-conditioned or ill-conditioned mathematical concepts (e.g. matrices in the most well-known context) arises. A procedure that is ill-conditioned applied to one problem might not be for another (e.g. inversion versus eigenvalue calculation in the case of matrices). Unfortunately it is in general impossible to determine at a glance whether it is the case or not and a detailed analysis might be required.

All effects must be understood and possibly minimized to extract the ”real physics” from the simulation. It is obvious that this should be a standard procedure for all types of simulations. It is often a very challenging task to identify possible problems. The omission of a careful analysis of these issues makes a simulation useless.

BENCHMARKING OF PROGRAMS

In the ideal world, all simulation results can be compared to reliable experimental data. To benchmark weak-strong and strong-strong beam-beam simulations is a challenge for the beam instrumentation.

Requirements for beam measurements

For a machine with large number of bunches the determination of parameters for individual bunches is almost inevitable. The (at least) qualitative agreement of simulations with single bunch measurements is a crucial test of the predictive power of a simulation program [1].

Although the simulation of the measurements can be (should be) part of the program, it is necessary to define a sensible set of observables that allow a definite verification of the results.

Furthermore, the good knowledge and control of all the other parameters of the machine such as optical parameters etc. is a prerequisite for a successful comparison.

Consistency between simulation programs

Last but not least it is necessary that programs simulating the same physics should get the same results. Since ”real physics” must not depend on the numerical procedures and algorithms, observations confirmed by independent programs increase significantly the confidence in the results. Although not explicitly mentioned during the panel discussion, we consider it very healthy and highly desirable that we have a set of tools available which use very different methods to extract the correct physics. This is in particular true in the field of strong-strong simulations which saw very significant leaps forward in the last few years. Development is still ongoing and promising.

REFERENCES

**SUMMARY OF THE PANEL DISCUSSION ON IMPEDANCE CODES**

T. Weiland, Institut TEMF, Technische Universität Darmstadt, D64285 Darmstadt, Germany

**Abstract**

During the CARE HHH workshop a panel discussion took place with the following members: K. Bane (SLAC), W. Bruns (Berlin), F. Caspers (CERN), M. Dohlus (DESY), E. Jensen (CERN, secretary), T. Weiland (TU-Darmstadt and CST, convenor) and M. Zobov (Frascati). The task of the panel was to review existing codes for impedance problems and identify areas where further development is needed.

**QUESTIONS**

All panellists have been asked to prepare answers for the following questions:

Is there any need for more software than already available today? And if so, what kind of?

(1) Do we need entirely new software packages or would it be sufficient to extend existing tools?

(2) Is there any completely new area of computation and type of physics which no software as of today can deal with?

(3) What are the pro's and con's of commercial and research software tools?

(4) What is the limitation when applying existing tools for current problems, memory or computational power?

(5) Do we need optimization tools in this area?

**ANSWERS**

All authors gave short presentations in which some tried to give answers to the questions above, some gave presentations on their experience in using the codes and some presented a specific impedance problem that they thought could be solved in the future. Table I summarizes the overview over the most important impedance codes.

The discussion following the short presentations gave a rather clear picture and common opinion that can be summarized as follows:

- There is no need for entirely new additional software tools.
- Both time and frequency domain tools are needed.
- Choice of meshing may be important.
- Optimization was not considered to be an urgent need in impedance computations.
- Instead of (re-)writing entirely new software specific add-ins to commercial software are preferred.
- Some extensions to resistive walls and thin sheets are still not covered well enough.

**REFERENCES**


<table>
<thead>
<tr>
<th>Tool</th>
<th>Origin</th>
<th>Method</th>
<th>Mesh</th>
<th>Specials</th>
</tr>
</thead>
<tbody>
<tr>
<td>CST-MICROWAVE STUDIO®</td>
<td>CST GmbH</td>
<td>FIT [2]</td>
<td>Choice of Tetrahedrons and</td>
<td>Commercial integrated optimizer network computing 32 Bit and 64 Bit versions</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Time Domain, Frequency Domain</td>
<td>PBA™ / Cartesian, Subgridding</td>
<td></td>
</tr>
<tr>
<td>ECHO 3D</td>
<td>TEMF TU Darmstadt</td>
<td>FIT [2]</td>
<td>PBA™ / Cartesian</td>
<td>Dedicated wake field codes, very long structures no grid dispersion</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Time Domain</td>
<td></td>
<td></td>
</tr>
<tr>
<td>GdfidL</td>
<td>W. Bruns</td>
<td>Same as MAFIA</td>
<td>Cartesian with Subcells</td>
<td>LINUX Cluster version, parallel version 32 Bit and 64 Bit versions</td>
</tr>
<tr>
<td>HFSS [4]</td>
<td>Ansoft</td>
<td>FE Frequency Domain</td>
<td>Tetrahedrons</td>
<td>Commercial, Frequency Domain only 32 Bit and 64 Bit versions</td>
</tr>
<tr>
<td>MAFIA [1]</td>
<td>CST GmbH</td>
<td>FIT [2]</td>
<td>Cartesian with Sub-cells; Sub-gridding</td>
<td>Complete package includes PIC 32 Bit and 64 Bit versions</td>
</tr>
<tr>
<td></td>
<td></td>
<td>Time Domain, Frequency Domain</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>Time Domain</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>
Abstract

I summarize the overview talk by M. Furman on ‘electron-cloud simulation codes’ and the subsequent electron-cloud panel discussion in Session 6 of the HHH-2004 workshop.

1 INTRODUCTION

The CARE-HHH-APD Workshop on ‘Beam Dynamics in Future Hadron Colliders and Rapidly Cycling High-Intensity Synchrotrons’ (HHH-2004) was held at CERN, November 8–11, 2004. Session 6 of the workshop was devoted to simulation code benchmarking and code repository, and consisted of two parts, addressing electron cloud and collective instabilities, respectively. Here, I summarize the electron-cloud discussions.

M. Furman of LBNL presented an overview of electron-cloud simulation codes. Afterwards a 60-minutes panel discussion was held. The panelists were

- A. Adelmann (PSI)
- G. Bellodi (CCLRC/Astec/Rutherford)
- M. Furman (LBNL)
- K. Ohmi (KEK)
- M. Pivi (SLAC)
- G. Rumolo (GSI)
- D. Schulte (CERN)
- J. Wei (BNL)
- F. Zimmermann (CERN)

The panel was asked to address the following questions:

- What information should future codes provide that cannot be obtained from today’s codes?
- Which missing physics should they include?
- Do we need entirely new software packages or would it be sufficient to extend existing tools?
- Would a common code repository be helpful to the community and, if yes, what features should it include?
- Which codes are you using and which improvements to these would you suggest?
- How can we quantify the reliability of the code predictions?
- Which kinds of benchmarking between codes and between simulations and experiments are needed or desirable?

2 GOALS

J. Wei emphasized that the codes should help address performance limiting issues, like

- vacuum pressure rise (RHIC),
- instabilities (PSR), and
- emittance growth (SPS, KEKB).

M. Furman stated that reliable predictions are the ultimate goal. In a similar spirit, D. Schulte expressed as the goal that the codes should predict conditions before an accelerator is built.

K. Ohmi proposed as a practical challenges the task to reproduce in simulations the KEKB observations, in particular the ‘sideband’ and relaxation oscillations seen in beam position and beam size (see Figs. 1 and 2), which are still an open issue, and the multibunch instability, which has successfully been reproduced in KEK simulations (Figs. 3 and 4).

3 TYPES OF CODES

M. Furman distinguished four types of electron-cloud simulation codes:
Figure 2: Time series data from (top) BPM, sensitive to the beam position, and (bottom) photomultiplier (PMT), sensitive to the beam size [1], revealing a sudden blow up and subsequent relaxation. Measurement from the KEKB LER.

Figure 3: Measured and simulated horizontal (top) and vertical (bottom) coupled bunch mode spectrum in the KEKB LER with solenoids off [2, 3].

Figure 4: Measured and simulated horizontal (top) and vertical (bottom) coupled bunch mode spectrum in the KEKB LER with solenoids on [2, 3].

Figure 5: Schematic of the ‘ultimate’ electron-cloud code.

1. electron-cloud build-up codes, where the beam is prescribed, except possibly for multibunch dipole motion, and the electrons modeled by macroparticles are dynamical; various chamber geometries, electron sources, and magnetic fields can be simulated;

2. instability codes, where the electron cloud is prescribed, at least initially, either as a (non-)linear lens or as a cloud of particles, and the beam is dynamical, represented by an ensemble of macroparticles;

3. self-consistent codes, which can exhibit various degrees of self-consistency, and whose main characteristic is to treat both beam and electrons dynamically; typically the self-consistent codes are 3D, and some accept an input lattice description; the self-consistent code may or may not model the secondary emission process at the chamber wall; ultimate goal of a self-consistent approach would be to model gas desorption, photo-electric effect, ionization, stray particles and particle-wall collisions, as well as secondary ionization;

4. map codes like MAC [4].

Table 1 lists frequently used electron-cloud codes with some of their properties.

4 OUTLOOK

Figure 5 shows a schematic of all the features which an ultimate electron-cloud code should either include or be interfaced to, and which may or may not interplay with the electron cloud, depending on the application considered.

5 ACKNOWLEDGEMENT

I thank M. Furman for his excellent overview talk and all electron-cloud panelists in Session 6 for the informative,
Table 1: Electron-cloud codes. Abbreviations are SC: self-consistent; SE=secondary electron emission, SR=synchrotron radiation, IZ=ionization of residual gas, and BPL=beam particle losses.

<table>
<thead>
<tr>
<th>code</th>
<th>contact</th>
<th>dim.</th>
<th>e− model</th>
<th>features</th>
<th>parallel (max. CPU)</th>
</tr>
</thead>
<tbody>
<tr>
<td>PEI</td>
<td>K. Ohmi, KEK</td>
<td></td>
<td>SR, SE</td>
<td>build-up dipole inst.</td>
<td></td>
</tr>
<tr>
<td>EPI</td>
<td>K. Ohmi, KEK</td>
<td></td>
<td>SR, SE</td>
<td>build up, dipole inst.</td>
<td></td>
</tr>
<tr>
<td>CLOUDLAND</td>
<td>L. Wang, BNL</td>
<td></td>
<td>SR, SE</td>
<td>build up</td>
<td>N</td>
</tr>
<tr>
<td>ECL CLOUD</td>
<td>G. Rumolo, GSI</td>
<td>2.5</td>
<td>SR, SE, IZ, BPL</td>
<td>build up, multi-bunch dipole wake</td>
<td>N</td>
</tr>
<tr>
<td>POS INST</td>
<td>M. Furman, LBNL</td>
<td>2.5</td>
<td>SE, IZ, BPL</td>
<td>build up, multipract dipole inst.</td>
<td></td>
</tr>
<tr>
<td>CSEC</td>
<td>M. Blaskiewicz, BNL</td>
<td>2–3</td>
<td>SE, IZ, BPL</td>
<td>build-up, single-bunch instability</td>
<td></td>
</tr>
<tr>
<td>HEAD TAIL</td>
<td>G. Rumolo, GSI</td>
<td>2</td>
<td>SE, IZ</td>
<td>build up single-bunch inst. (PIC)</td>
<td></td>
</tr>
<tr>
<td>PEHT</td>
<td>K. Ohmi, KEK</td>
<td></td>
<td></td>
<td>head-tail (microbunches)</td>
<td></td>
</tr>
<tr>
<td>PEHT S</td>
<td>K. Ohmi, KEK</td>
<td></td>
<td></td>
<td>head-tail, SC (PIC)</td>
<td></td>
</tr>
<tr>
<td>CLOUD MAD</td>
<td>T. Raubenheimer, SLAC</td>
<td></td>
<td></td>
<td>MAD tracking with e− cloud lenses</td>
<td></td>
</tr>
<tr>
<td>PARSEC</td>
<td>A. Adelmann, PSI</td>
<td>3</td>
<td>SE, IZ, SR, BPL</td>
<td>SC; lattice description</td>
<td>Y (4048)</td>
</tr>
<tr>
<td>ORBIT</td>
<td>J. Holmes, ORNL</td>
<td>2–3</td>
<td>SE, IZ</td>
<td>SC; lattice description</td>
<td>Y</td>
</tr>
<tr>
<td>WARP++POS INST</td>
<td>J.L. Vay, LBNL</td>
<td>4</td>
<td>SE, IZ, SR, BPL</td>
<td>SC, lattice description</td>
<td>Y</td>
</tr>
<tr>
<td>QUICKPIC</td>
<td>T. Katsoulas, USC</td>
<td>2–3</td>
<td>PIC plasma code; initially prescribed cloud</td>
<td>Y (128)</td>
<td></td>
</tr>
<tr>
<td>BEST</td>
<td>H. Win, PPFL</td>
<td>3</td>
<td>SC, Vlasov-Maxwell, no e-wall collisions</td>
<td>Y (512)</td>
<td></td>
</tr>
<tr>
<td>MEC</td>
<td>U. Iriso, BNL</td>
<td></td>
<td></td>
<td>empirical maps</td>
<td></td>
</tr>
</tbody>
</table>
well prepared and highly inspiring discussion. I also thank F. Ruggiero for some comments on the summary.

6 REFERENCES


Abstract

This report is a brief summary of the session 5 and the panel discussion 6 of the workshop CARE-HHH, which were devoted to coherent instabilities as potential mechanisms of intensity limitations in future machines. In this session comprehensive overviews of the present knowledge on collective phenomena as well as on frontier research for novel aspects were presented. The status and limitations of some of the existing (analytical and numerical) methods to study transverse and longitudinal collective effects have been reported and discussed by the panel members. Questions about validity, benchmarking and desirable computational speed of codes for large scale simulations (space charge, impedances) have been raised and, where possible, addressed.

LIST OF THE CONTRIBUTIONS

Session 5 consisted of 2 invited talks and a panel discussion:

- **Overview of single-beam conventional collective instabilities in both transverse and longitudinal planes**, Elias Métral (CERN)
- **Intensity limitations by combined and/or (un)conventional impedance sources**, Giovanni Rumolo (GSI/Università di Napoli “Federico II”)
- **Panel Discussion 6: Coherent Instabilities**, A. Adelmann (PSI), E. Métral (CERN), L. Palumbo (INFN-LNF), M. Zobov (INFN-LNF)

TALK SUMMARIES

**Overview of single-beam conventional collective instabilities in both transverse and longitudinal planes**

Collective instabilities can in principle occur at whatever beam intensity because of the wrong sign of chromaticity or because of the operating conditions which can make the threshold very low. At low intensity the different modes of oscillation (head-tail for the transverse plane and longitudinal) are standing-wave patterns, which can be treated independently (Sacherer formalism). In particular longitudinal coherent frequency of the dipole mode moves inside the incoherent frequency spectrum as the synchrotron frequency spread increases. Experimental evidence of the existence of these modes in the transverse plane was for example observed at the CERN-PS in 1999, where the harmonic number on which the instability arose (6) could be predicted beforehand via an analytical model that showed 6 as the first (and quickest) unstable mode.

At high intensity the modes can no longer be treated separately and mode-coupling between the different modes has to be taken into consideration.

- For the transverse plane, 5 different formalisms (coasting beam approach with peak current values, fast blow up theory by Ruth and Wang, beam break up from linear accelerators, post head-tail from Kernel et al., transverse mode coupling formalism with 2 modes in the long bunch regime) give the same threshold value for instability within a factor 2. Transverse Mode Coupling Instability (TMCI), or Strong Head-tail Instability, has been observed for the first time in a proton machine at the CERN-SPS in 2003. The threshold from the coasting beam formula with peak current values has been compared to thresholds from the MOSES code (by Chin, it is a computer program solving the mode-coupling equation and giving the threshold current at which mode coupling occurs first) and from the HEADTAIL code (by G. Rumolo, it is a macroparticle code for transverse single-bunch phenomena) and results of the comparisons (code-analytical formula, code-code, code-measurements) have been presented at the ICFA-HB2004 Workshop, held in Bensheim (Germany), 18-22 October 2004. The threshold can be raised by increasing the longitudinal emittance or the chromaticity.

- For the longitudinal plane, a formula is given, which comes from mode-coupling between the 2 most critical modes and takes into account the potential-well distortion. The use of this formula makes clear that: 1) below transition the beam can be stable even with a space-charge impedance much larger than the broadband impedance, and 2) it is better to inject into a machine below transition than above.

Finally the stabilizing methods for low-intensity cases were discussed:

- Landau damping (for the transverse plane)
  - From octupoles only (S. Berg and F. Ruggiero).
  - From both octupoles and space-charge nonlinearities (Mohl-Schonauer prediction of loss of Landau damping is recovered).
- Next work: include the longitudinal motion.
- Active feedbacks.
- Linear coupling between the transverse planes (with and without external, i.e. from octupoles, nonlinear-
In particular this method is used in the CERN PS machine to stabilize the beam for LHC.

**Intensity limitations by combined and/or (un)conventional impedance sources**

Impedance sources can be classified as "conventional", which describe the interaction of the beam with itself and with the “geometrically” surrounding environment (like space charge, resistive wall, broad- or narrow-band resonators), or “unconventional”, which describe the interaction of the beam with “external” fields induced by a medium, like an electron plasma, or radiated by the beam itself (coherent synchrotron radiation, electron cooler, electron cloud). The features of the wake fields associated to unconventional impedance sources are generally different and require a more complex description than those of conventional wake fields. It is interesting to study both the effect of an unconventional wake field on a conventional bunch and the effect of a conventional impedance (e.g., a broad-band impedance) on an unconventional bunch (e.g. a bunch in a barrier bucket)

Electron cloud wake fields and impedances have been first introduced in 1999 by P. Zenkevich for coasting beams, considering only electrons generated by residual gas ionization and trapped around the beam. For bunched beams the question required an extension of the concept, since electrons are not only primarily generated by bunches through residual gas ionization, but they also multipact due to acceleration in the beam field and secondary emission at the pipe walls. The unconventional features of electron cloud wake fields can be summarized as follows:

- The dipole wake field of an electron cloud depends on the transverse coordinates \((x, y)\)
- Differently located displacements along a bunch create differently shaped wake fields
- The wake field depends:
  - Strongly:
    - On the initial electron distribution
    - On the bunch particle transverse distribution
  - Weakly:
    - On the boundary conditions for a wide pipe
    - On the electron space charge for low degrees of neutralization
- A description in terms of double frequency impedance \(Z(\omega, \omega')\) is necessary for a correct Transverse Mode Coupling (TMC) analysis

Numerical tools to handle the calculation of \(Z(\omega, \omega')\) have been developed, but both the dependence of the wake on the longitudinal shape of the bunch and the electron cloud wake field for long bunches (which might strongly depend on the trailing edge electron production and multiplication) are yet to be investigated.

An accurate analysis of dipole and envelope modes of a bunch in a barrier bucket led to the following conclusions:

- The coherent tune shift \(\Delta Q\) of a bunch in a barrier bucket as a function of the bunch current depends on
  - Shunt impedance (proportional)
  - Bunch length (inversely proportional) and maybe momentum spread (proportional?)
  - Chamber shape (only in \(x\))
- The \(\Delta Q\) in low current follows that of a usual bunch in a sinusoidal bucket with the same longitudinal emittance (theoretical line).
- Coherent envelope modes depend on the chamber shape:
  - Round chamber has two modes both in \(x\) and \(y\), one current dependent and one current independent.
  - Flat chamber has one mode in \(x\) with a positive shift with increasing current, and two modes in \(y\), both with a negative shift with current.

**Concerning instabilities:**

- High current: The threshold for strong head-tail instability is not found for bunches in a barrier bucket, but there is rather a regime of slow growth at high currents.
- Current independent: Regular head-tail instability driven by negative \(Q^*\) (above transition) exhibits similar features as for bunches in sinusoidal buckets.
  - Growth rates are proportional to the shunt impedance
  - The quickest instability occurs when \(\omega_\xi = \omega_r\).
  - In a flat chamber growth times in the \(x\) direction are about double of the growth times in the \(y\) direction
  - Longer bunches slow down the instability

An analytical model (maybe few particles model or kinetic model based on Vlasov equation) is needed.

**PANEL DISCUSSION**

The following questions were raised by the chairman to start the discussion between the panel members:

- Which physics do present simulation codes include and which physics should be included?
- Which physics would deserve more analytical studies?
- Code benchmark
- What are the code limitations?

L. Palumbo argued that the method usually employed to calculate wake fields in a variety of structures (pseudo-Green function, i.e. the electromagnetic response of the structure to the passage of a very thin disc or a ring of charge) can introduce an error in the high frequency range because of the artificial decaying spectrum of the source above some frequency (a point-like charge would have instead a uniform spectrum all over the frequency axis).
Some micro-bunching phenomena could then be not correctly taken into account when using wake fields from pseudo-Green functions. Referring in particular to the longitudinal plane, L. Palumbo listed a few tools used to study instability thresholds and unstable dynamics (Haissinski equation solvers, mode coupling theory from Vlasov equation, tracking codes), and said they have mostly proven reliable, even if there is a lack of: 1) analytical description of the beam evolution above the instability threshold, and 2) of systematic benchmarking of different approaches and codes.

A. Adelmann mainly devoted his statement to describing the present efforts to set up massive 3D and parallel space charge calculations for beam dynamics. More and more powerful parallel clusters are presently starting operation and new algorithms are being developed (like Finite Element Method on MultiGrid), thanks to which accurate and optimized macro-particle simulations will be carried out, which can include a description of space charge and interaction with the environment to whatever desired degree of detail.

E. Métral pointed out that lately a quite accurate benchmark between MOSES (code using the transverse mode coupling calculation from the kinetic theory) and HEADTAIL (macroparticle tracking code for single bunch phenomena) was carried out using the study of the TMCI threshold in the CERN-SPS. The agreement between the predictions of the two codes is very good. The advantage of HEADTAIL with respect to MOSES is that it can easily include in the simulation more realistic features like flat beam chamber, space charge, sinusoidal bucket, possibly unmatched bunch executing quadrupole oscillations.

M. Zobov extended L. Palumbo’s previous remark about the lack of analytical theory above the instability threshold. Also in the transverse plane, there is no satisfactory theory to describe the beam dynamics in unstable regime. M. Zobov also showed some potentially serious impedance sources for LHC (collimators and scrapers) and some instability observations in the positron ring of DaΦne, suggesting that it could be induced by electron cloud alone or by a combined effect resistive wall-electron cloud. K. Ohmi showed plots of the long-range wake field of an electron cloud, which could indeed be responsible for the observed instability.

The two following questions, which were indicated as potential concerns for the performance of LHC, arose:

- Has the longitudinal impedance of collimators in LHC been taken into account, and have simulations of longitudinal dynamics been done in order to make sure that it will not cause major troubles?
- Could space charge play a significant role even at top energy, since the storage time is very long?

**ACKNOWLEDGMENTS**

First of all, we are sincerely grateful to the CARE-HHH organizers F. Ruggiero, W. Scandale and F. Zimmermann for inviting us to give an active contribution to the success of the workshop. We also would like to thank all the speakers and panel members. They have done an excellent job creating an extremely rich and interesting session. Besides, lots of post-session fruitful discussion and plans for collaboration were possible thanks to E. Shaposhnikova, G. Franchetti and F. Zimmermann. Finally, many thanks to J. Thomashausen, A. Mostacci and E. Métral for the great help and practical support they gave us in preparing the summary talk on the last day of the workshop.
### List of HHH-2004 Participants

<table>
<thead>
<tr>
<th>Name</th>
<th>Affiliation</th>
</tr>
</thead>
<tbody>
<tr>
<td>Adelmann</td>
<td>Andreas GSI &amp; LBNL</td>
</tr>
<tr>
<td>Aleksan</td>
<td>Roy CEA Saclay</td>
</tr>
<tr>
<td>Altarelli</td>
<td>Guido CERN</td>
</tr>
<tr>
<td>Assmann</td>
<td>Ralph CERN</td>
</tr>
<tr>
<td>Aymar</td>
<td>Robert CERN</td>
</tr>
<tr>
<td>Baglin</td>
<td>Bagline CERN</td>
</tr>
<tr>
<td>Bane</td>
<td>Karl SLAC</td>
</tr>
<tr>
<td>Bellodi</td>
<td>Giulia Rutherford</td>
</tr>
<tr>
<td>Benedetto</td>
<td>Elena CERN</td>
</tr>
<tr>
<td>Benedikt</td>
<td>Michael CERN</td>
</tr>
<tr>
<td>Bordry</td>
<td>Frederick CERN</td>
</tr>
<tr>
<td>Bottura</td>
<td>Luca CERN</td>
</tr>
<tr>
<td>Blondel</td>
<td>Alain CERN</td>
</tr>
<tr>
<td>Bruning</td>
<td>Oliver CERN</td>
</tr>
<tr>
<td>Bruns</td>
<td>Warner CERN (Warner Bruns Feldberechnungen)</td>
</tr>
<tr>
<td>Caspers</td>
<td>Fritz CERN</td>
</tr>
<tr>
<td>Cervelli</td>
<td>Franco INFN</td>
</tr>
<tr>
<td>Chiaveri</td>
<td>Enrico CERN</td>
</tr>
<tr>
<td>Cimino</td>
<td>Roberto INFN Frascati</td>
</tr>
<tr>
<td>Collier</td>
<td>Paul CERN</td>
</tr>
<tr>
<td>Damerau</td>
<td>Heiko CERN</td>
</tr>
<tr>
<td>Debu</td>
<td>Pascal CEA Saclay</td>
</tr>
<tr>
<td>Decking</td>
<td>Winfried DESY</td>
</tr>
<tr>
<td>Delahaye</td>
<td>Jean-Pierre CERN</td>
</tr>
<tr>
<td>Demma</td>
<td>Theo Univ. of Sannio</td>
</tr>
<tr>
<td>Denegri</td>
<td>Daniel CERN</td>
</tr>
<tr>
<td>Dohlus</td>
<td>Martin DESY</td>
</tr>
<tr>
<td>Dosselli</td>
<td>Umberto INFN Padova</td>
</tr>
<tr>
<td>Engelen</td>
<td>Jos CERN</td>
</tr>
<tr>
<td>Faus-Golfe</td>
<td>Angeles Instituto de Fisica Corpuscular, Valencia</td>
</tr>
<tr>
<td>Ferrini</td>
<td>Federico Permanent Mission of Italy</td>
</tr>
<tr>
<td>Ferrari</td>
<td>Alfredo CERN</td>
</tr>
<tr>
<td>Formenti</td>
<td>Fabio CERN</td>
</tr>
<tr>
<td>Franchetti</td>
<td>Giuliano GSI</td>
</tr>
<tr>
<td>Furman</td>
<td>Miguel LBNL</td>
</tr>
<tr>
<td>Garoby</td>
<td>Roland CERN</td>
</tr>
<tr>
<td>Giovannozzi</td>
<td>Massimo CERN</td>
</tr>
<tr>
<td>Goddard</td>
<td>Brennan CERN</td>
</tr>
<tr>
<td>Gourlay</td>
<td>Steves LBNL</td>
</tr>
<tr>
<td>Gröbner</td>
<td>Oswald CERN</td>
</tr>
<tr>
<td>Guidi</td>
<td>Vincenzo Università di Ferrara</td>
</tr>
<tr>
<td>Henning</td>
<td>Walter F. GSI</td>
</tr>
<tr>
<td>Herr</td>
<td>Werner CERN</td>
</tr>
<tr>
<td>Jeanneret</td>
<td>Jean-Bernard CERN</td>
</tr>
<tr>
<td>Jensen</td>
<td>Erk CERN</td>
</tr>
<tr>
<td>Jimenez</td>
<td>Jose Miguel CERN</td>
</tr>
<tr>
<td>Koutchouk</td>
<td>Jean-Pierre CERN</td>
</tr>
<tr>
<td>Kroyer</td>
<td>Tom CERN</td>
</tr>
</tbody>
</table>

